

# A Two-Singlet Model for Light Cold Dark Matter

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## Abstract

We extend the Standard Model by adding two gauge-singlet  $\mathbb{Z}_2$ -symmetric scalar fields that interact with visible matter only through the Higgs particle. One is a stable dark matter WIMP, and the other one undergoes a spontaneous breaking of the symmetry that opens new channels for the dark matter annihilation, hence lowering the mass of the WIMP. We study the effects of the observed dark matter relic abundance on the annihilation cross section and find that in most regions of the parameters space, light dark matter is viable. We also compare the elastic scattering cross-section of our dark matter candidate off nucleus with existing (CDMSII and XENON100) and projected (SuperCDMS and XENON1T) experimental exclusion bounds. We find that most of the allowed mass range for light dark matter will be probed by the projected sensitivity of XENON1T experiment.

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## I. INTRODUCTION

Cosmology tells us that about 25% of the total mass density in the Universe is dark matter that cannot be accounted for by conventional baryons [1]. Alongside observation, intense theoretical efforts are made in order to elucidate the nature and properties of this unknown form of matter. In this context, electrically neutral and colorless weakly interacting massive particles (WIMPs) form an attractive scenario. Their broad properties are: masses in the range of one to a few hundred GeV, coupling constants in the milli-weak scale and lifetimes longer than the age of the Universe.

Recent data from the direct-detection experiments DAMA/LIBRA [2] and CoGeNT [3], and the recent analysis of the data from the Fermi Gamma Ray Space Telescope [4], if interpreted as signal for dark matter, require light WIMPs in the range of 5 to 10 GeV [5]. Also, galactic substructure requires still lighter dark matter masses [6, 7]. In this regard, it is useful to note in passing that the XENON100 collaboration has provided serious constraints on the region of interest to DAMA/LIBRA and CoGeNT [8], assuming a constant extrapolation of the liquid xenon scintillation response for nuclear recoils below 5 keV, a claim disputed in [9]. Most recently, the CDMS collaboration has released the analysis of their low-energy threshold data [10] which seems to exclude the parameter space for dark matter interpretation of DAMA/LIBRA and CoGeNT results, assuming a standard halo dark matter model with an escape velocity  $v_{esc} = 544$  km/s and neglecting the effect of ion channeling [11]. However, with a highly anisotropic velocity distribution, it may be possible to reconcile the CoGeNT and DAMA/LIBRA results with the current exclusion limits from CDMS and XENON [12] (see also comments on p.6 in [13] about the possibility of shifting the exclusion contour in [10] above the CoGeNT signal region). In addition, CRESST, another direct detection experiment at Gran Sasso, which uses CaWO<sub>4</sub> as target material, reported in talks at the IDM 2010 and WONDER 2010 workshops an excess of events in their oxygen band instead of tungsten band. If this signal is not due to neutron background a possible interpretation could be the elastic scattering of a light WIMP depositing a detectable recoil energy on the the lightest nuclei (oxygen) in the detector [14]. While this result has to await confirmation from the CRESST collaboration, it is clear that it is important as well as interesting to study light dark matter.

The most popular candidate for dark matter is the neutralino, a neutral  $R$ -odd supersymmetric particle. Indeed, they are only produced or destroyed in pairs, thus rendering the lightest SUSY particle (LSP) stable [15]. In the minimal version of the supersymmetric extension of the Standard Model, the neutralino is a linear combination of the fermionic partners of the neutral electroweak gauge bosons (gauginos) and the neutral Higgs bosons (higgsinos). They can annihilate through a t-channel sfermion exchange into standard model fermions, or via a t-channel chargino-mediated process into  $W^+W^-$ , or through an s-channel pseudoscalar Higgs exchange into fermion pairs and they can undergo elastic scattering with nuclei mainly through scalar Higgs exchange [16]. If these

neutralinos were light, they would then be overproduced in the early universe and, in the minimal model, would not have an elastic-scattering cross section large enough to account for the DAMA/LIBRA and CoGeNT results due to constraints from other experiments such as the LEP, Tevatron, and rare decays [17–19].

Therefore, with no clear clue yet as to what the internal structure of these WIMPs is, if any, a pedestrian approach can be attractive. In this logic, the simplest of models is to extend the Standard Model by adding a real scalar field, the dark matter, a Standard-Model gauge singlet that interacts with visible particles via the Higgs field only. To ensure stability, it is endowed with a discrete  $\mathbb{Z}_2$  symmetry that does not spontaneously break. Such a model can be seen as a low-energy remnant of some higher-energy physics waiting to be understood. In this cosmological setting, such an extension has first been proposed in [20] and further studied in [21] where the unbroken  $\mathbb{Z}_2$  symmetry is extended to a global U(1) symmetry. A more extensive exploration of the model and its implications was done in [22], specific implications on Higgs detection and LHC physics discussed in [23] and one-loop vacuum stability looked into and perturbativity bounds obtained in [24]. The work of [25] considers also this minimal extension and uses constraints from the experiments XENON10 [26] and CDMSII [27] to exclude dark matter masses smaller than 50, 70 and 75GeV for Higgs masses equal to 120, 200 and 350 GeV respectively.

In order to allow for light dark matter, it is therefore necessary to go beyond the minimal one-real-scalar extension of the Standard Model. The natural next step is to add another real scalar field, endowed with a  $\mathbb{Z}_2$  symmetry too, but one which is spontaneously broken so that new channels for dark matter annihilation are opened, increasing this way the annihilation cross-section, hence allowing smaller masses. This auxiliary field must also be a Standard Model gauge singlet.

After this brief introductory motivation, we present the model in the next section. We perform the spontaneous breaking of the electroweak and the additional  $\mathbb{Z}_2$  symmetries in the usual way. We clarify the physical modes as well as the physical parameters. There is mixing between the physical new scalar field and the Higgs, and this is one of the quantities parametrizing the subsequent physics. In section three, we impose the constraint from the known dark matter relic density on the dark-matter annihilation cross section and study its effects. Of course, as we will see, the parameter space is quite large, and so, it is not realistic to hope to cover all of it in one single work of acceptable size. Representative values have to be selected and the behavior of the model as well as its capabilities are described. Our main focus in this study is the mass range 0.1GeV – 100GeV and we find that the model is rich enough to bear dark matter in most of it, including the very light sector. In section four, we determine the total cross section  $\sigma_{\text{det}}$  for non relativistic elastic scattering of dark matter off a nucleon target and compare it to the current direct-detection experimental bounds and projected sensitivity. For this, we choose the results of CDMSII and XENON100, and the projections of SuperCDMS [28] and XENON1T [29]. Here too we cannot cover all of the parameter space nor are we going to give a detailed account of the behavior of  $\sigma_{\text{det}}$  as a function of the dark matter

mass, but general patterns are mentioned. The last section is devoted to some concluding remarks. Note that as a rule, we have avoided in this first study narrowing the choice of parameters using particle phenomenology. Of course, such phenomenological constraints have to be addressed ultimately and this is left to a forthcoming investigation, contenting ourselves in the present work with a limited set of remarks mentioned in this last section. Finally, we have gathered in an appendix the partial results regarding the calculation of the dark matter annihilation cross section.

## II. A TWO-SINGLET MODEL FOR DARK MATTER

We extend the Standard Model by adding two real, spinless and  $\mathbb{Z}_2$ -symmetric fields: the dark matter field  $S_0$  for which the  $\mathbb{Z}_2$  symmetry is unbroken and an auxiliary field  $\chi_1$  for which it is spontaneously broken. Both fields are Standard Model gauge singlets and hence can interact with ‘visible’ particles only via the Higgs doublet  $H$ . This latter is taken in the unitary gauge such that  $H^\dagger = 1/\sqrt{2}(0 \ h')$ , where  $h'$  is a real scalar. We assume all processes calculable in perturbation theory. The potential function that includes  $S_0$ ,  $h'$  and  $\chi_1$  writes as follows:

$$U = \frac{\tilde{m}_0^2}{2}S_0^2 - \frac{\mu^2}{2}h'^2 - \frac{\mu_1^2}{2}\chi_1^2 + \frac{\eta_0}{24}S_0^4 + \frac{\lambda}{24}h'^4 + \frac{\eta_1}{24}\chi_1^4 + \frac{\lambda_0}{4}S_0^2h'^2 + \frac{\eta_{01}}{4}S_0^2\chi_1^2 + \frac{\lambda_1}{4}h'^2\chi_1^2, \quad (2.1)$$

where  $\tilde{m}_0^2$ ,  $\mu^2$  and  $\mu_1^2$  and all the coupling constants are real positive numbers. In the Standard Model scenario, electroweak spontaneous symmetry breaking occurs for the Higgs field, which then oscillates around the vacuum expectation value  $v = 246\text{GeV}$  [30]. The field  $\chi_1$  will oscillate around the vacuum expectation value  $v_1 > 0$ . Both  $v$  and  $v_1$  are related to the parameters of the theory by the two relations:

$$v^2 = 6\frac{\mu^2\eta_1 - 6\mu_1^2\lambda_1}{\lambda\eta_1 - 36\lambda_1^2}, \quad v_1^2 = 6\frac{\mu_1^2\lambda - 6\mu^2\lambda_1}{\lambda\eta_1 - 36\lambda_1^2}. \quad (2.2)$$

It is assumed that the self-coupling constants are sufficiently larger than the mutual ones.

Writing  $h' = v + \tilde{h}$  and  $\chi_1 = v_1 + \tilde{\chi}_1$ , the potential function becomes, up to an irrelevant zero-field energy:

$$U = U_{\text{quad}} + U_{\text{cub}} + U_{\text{quar}}, \quad (2.3)$$

where the mass-squared (quadratic) terms are gathered in  $U_{\text{quad}}$ , the cubic interactions in  $U_{\text{cub}}$  and the quartic ones in  $U_{\text{quar}}$ . The quadratic terms are given by:

$$U_{\text{quad}} = \frac{1}{2}m_0^2S_0^2 + \frac{1}{2}M_h^2\tilde{h}^2 + \frac{1}{2}M_1^2\tilde{\chi}_1^2 + M_{1h}^2\tilde{h}\tilde{\chi}_1, \quad (2.4)$$

where the mass-squared coefficients are related to the original parameters of the theory by the following relations:

$$\begin{aligned} m_0^2 &= \tilde{m}_0^2 + \frac{\lambda_0}{2}v^2 + \frac{\eta_{01}}{2}v_1^2; & M_h^2 &= -\mu^2 + \frac{\lambda}{2}v^2 + \frac{\lambda_1}{2}v_1^2; \\ M_1^2 &= -\mu_1^2 + \frac{\lambda_1}{2}v^2 + \frac{\eta_1}{2}v_1^2; & M_{1h}^2 &= \lambda_1 v v_1. \end{aligned} \quad (2.5)$$

Replacing the vacuum expectation values  $v$  and  $v_1$  by their respective expressions (2.2) will not add clarity. In this field basis, the mass-squared matrix is not diagonal: there is mixing between the fields  $\tilde{h}$  and  $\tilde{S}_1$ . Denoting the physical mass-squared field eigenmodes by  $h$  and  $S_1$ , we rewrite:

$$U_{\text{quad}} = \frac{1}{2}m_0^2 S_0^2 + \frac{1}{2}m_h^2 h^2 + \frac{1}{2}m_1^2 S_1^2, \quad (2.6)$$

where the physical fields are related to the mixed ones by a  $2 \times 2$  rotation:

$$\begin{pmatrix} h \\ S_1 \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} \tilde{h} \\ \tilde{S}_1 \end{pmatrix}. \quad (2.7)$$

Here  $\theta$  is the mixing angle, related to the original mass-squared parameters by the relation:

$$\tan 2\theta = \frac{2M_{1h}^2}{M_1^2 - M_h^2}, \quad (2.8)$$

and the physical masses in (2.6) by the two relations:

$$\begin{aligned} m_h^2 &= \frac{1}{2} \left[ M_h^2 + M_1^2 + \varepsilon (M_h^2 - M_1^2) \sqrt{(M_h^2 - M_1^2)^2 + 4M_{1h}^4} \right]; \\ m_1^2 &= \frac{1}{2} \left[ M_h^2 + M_1^2 - \varepsilon (M_h^2 - M_1^2) \sqrt{(M_h^2 - M_1^2)^2 + 4M_{1h}^4} \right], \end{aligned} \quad (2.9)$$

where  $\varepsilon$  is the sign function.

Written now directly in terms of the physical fields, the cubic interaction terms are expressed as follows:

$$U_{\text{cub}} = \frac{\lambda_0^{(3)}}{2} S_0^2 h + \frac{\eta_{01}^{(3)}}{2} S_0^2 S_1 + \frac{\lambda^{(3)}}{6} h^3 + \frac{\eta_1^{(3)}}{6} S_1^3 + \frac{\lambda_1^{(3)}}{2} h^2 S_1 + \frac{\lambda_2^{(3)}}{2} h S_1^2, \quad (2.10)$$

where the cubic physical coupling constants are related to the original parameters via the following relations:

$$\begin{aligned} \lambda_0^{(3)} &= \lambda_0 v \cos \theta + \eta_{01} v_1 \sin \theta; \\ \eta_{01}^{(3)} &= \eta_{01} v_1 \cos \theta - \lambda_0 v \sin \theta; \\ \lambda^{(3)} &= \lambda v \cos^3 \theta + \frac{3}{2} \lambda_1 \sin 2\theta (v_1 \cos \theta + v \sin \theta) + \eta_1 v_1 \sin^3 \theta; \\ \eta_1^{(3)} &= \eta_1 v_1 \cos^3 \theta - \frac{3}{2} \lambda_1 \sin 2\theta (v \cos \theta - v_1 \sin \theta) - \lambda v \sin^3 \theta; \\ \lambda_1^{(3)} &= \lambda_1 v_1 \cos^3 \theta + \frac{1}{2} \sin 2\theta [(2\lambda_1 - \lambda) v \cos \theta - (2\lambda_1 - \eta_1) v_1 \sin \theta] - \lambda_1 v \sin^3 \theta; \\ \lambda_2^{(3)} &= \lambda_1 v \cos^3 \theta - \frac{1}{2} \sin 2\theta [(2\lambda_1 - \eta_1) v_1 \cos \theta + (2\lambda_1 - \lambda) v \sin \theta] + \lambda_1 v_1 \sin^3 \theta. \end{aligned} \quad (2.11)$$

Also, in terms of the physical fields, the quartic interactions are given by:

$$\begin{aligned} U_{\text{quar}} = & \frac{\eta_0}{24} S_0^4 + \frac{\lambda^{(4)}}{24} h^4 + \frac{\eta_1^{(4)}}{24} S_1^4 + \frac{\lambda_0^{(4)}}{4} S_0^2 h^2 + \frac{\eta_{01}^{(4)}}{4} S_0^2 S_1^2 + \frac{\lambda_{01}^{(4)}}{2} S_0^2 h S_1 \\ & + \frac{\lambda_1^{(4)}}{6} h^3 S_1 + \frac{\lambda_2^{(4)}}{4} h^2 S_1^2 + \frac{\lambda_3^{(4)}}{6} h S_1^3, \end{aligned} \quad (2.12)$$

where the physical quartic coupling constants are written in terms of the original parameters of the theory as follows:

$$\begin{aligned} \lambda^{(4)} &= \lambda \cos^4 \theta + \frac{3}{2} \lambda_1 \sin^2 2\theta + \eta_1 \sin^4 \theta; \\ \eta_1^{(4)} &= \eta_1 \cos^4 \theta + \frac{3}{2} \lambda_1 \sin^2 2\theta + \lambda \sin^4 \theta; \\ \lambda_0^{(4)} &= \lambda_0 \cos^2 \theta + \eta_{01} \sin^2 \theta; \\ \eta_{01}^{(4)} &= \eta_{01} \cos^2 \theta + \lambda_0 \sin^2 \theta; \\ \lambda_{01}^{(4)} &= \frac{1}{2} (\eta_{01} - \lambda_0) \sin 2\theta; \\ \lambda_1^{(4)} &= \frac{1}{2} [(3\lambda_1 - \lambda) \cos^2 \theta - (3\lambda_1 - \eta_1) \sin^2 \theta] \sin 2\theta; \\ \lambda_2^{(4)} &= \lambda_1 \cos^2 2\theta - \frac{1}{4} (2\lambda_1 - \eta_1 - \lambda) \sin^2 2\theta; \\ \lambda_3^{(4)} &= \frac{1}{2} [(\eta_1 - 3\lambda_1) \cos^2 \theta - (\lambda - 3\lambda_1) \sin^2 \theta] \sin 2\theta. \end{aligned} \quad (2.13)$$

Finally, after spontaneous breaking of the electroweak and  $\mathbb{Z}_2$  symmetries, the part of the Standard Model lagrangian that is relevant to dark matter annihilation writes, in terms of the physical fields  $h$  and  $S_1$ , as follows:

$$\begin{aligned} U_{\text{SM}} = & \sum_f (\lambda_{hf} h \bar{f} f + \lambda_{1f} S_1 \bar{f} f) + \lambda_{hw}^{(3)} h W_\mu^- W^{+\mu} + \lambda_{1w}^{(3)} S_1 W_\mu^- W^{+\mu} \\ & + \lambda_{hz}^{(3)} h (Z_\mu)^2 + \lambda_{1z}^{(3)} S_1 (Z_\mu)^2 + \lambda_{hw}^{(4)} h^2 W_\mu^- W^{+\mu} + \lambda_{1w}^{(4)} S_1^2 W_\mu^- W^{+\mu} \\ & + \lambda_{h1w} h S_1 W_\mu^- W^{+\mu} + \lambda_{hz}^{(4)} h^2 (Z_\mu)^2 + \lambda_{1z}^{(4)} S_1^2 (Z_\mu)^2 + \lambda_{h1z} h S_1 (Z_\mu)^2. \end{aligned} \quad (2.14)$$

The quantities  $m_f$ ,  $m_w$  and  $m_z$  are the masses of the fermion  $f$ , the  $W$  and the  $Z$  gauge bosons respectively, and the above coupling constants are given by the following relations:

$$\begin{aligned} \lambda_{hf} &= -\frac{m_f}{v} \cos \theta; & \lambda_{1f} &= \frac{m_f}{v} \sin \theta; \\ \lambda_{hw}^{(3)} &= 2 \frac{m_w^2}{v} \cos \theta; & \lambda_{1w}^{(3)} &= -2 \frac{m_w^2}{v} \sin \theta; \\ \lambda_{hz}^{(3)} &= \frac{m_z^2}{v} \cos \theta; & \lambda_{1z}^{(3)} &= -\frac{m_z^2}{v} \sin \theta; \\ \lambda_{hw}^{(4)} &= \frac{m_w^2}{v^2} \cos^2 \theta; & \lambda_{1w}^{(4)} &= \frac{m_w^2}{v^2} \sin^2 \theta; & \lambda_{h1w} &= -\frac{m_w^2}{v^2} \sin 2\theta; \\ \lambda_{hz}^{(4)} &= \frac{m_z^2}{2v^2} \cos^2 \theta; & \lambda_{1z}^{(4)} &= \frac{m_z^2}{2v^2} \sin^2 \theta; & \lambda_{h1z} &= -\frac{m_z^2}{2v^2} \sin 2\theta. \end{aligned} \quad (2.15)$$

### III. RELIC DENSITY, MUTUAL COUPLINGS AND PERTURBATIVITY

The original theory (2.1) has nine parameters: three mass parameters  $(\tilde{m}_0, \mu, \mu_1)$ , three self-coupling constants  $(\eta_0, \lambda, \eta_1)$  and three mutual coupling constants  $(\lambda_0, \eta_{01}, \lambda_1)$ . Perturbativity is assumed, hence all these original coupling constants are small. The dark-matter self-coupling constant  $\eta_0$  does not enter in the calculations of the lowest-order processes of this work [31], so effectively, we are left with eight parameters. The spontaneous breaking of the electroweak and  $\mathbb{Z}_2$  symmetries for the Higgs and  $\chi_1$  fields respectively introduces the two vacuum expectation values  $v$  and  $v_1$  given to lowest order in (2.2). The value of  $v$  is fixed experimentally to be 246GeV and for the present work, we fix the value of  $v_1$  at the order of the electroweak scale, say 100GeV. Hence we are left with six parameters. Four of these are chosen to be the three physical masses  $m_0$  (dark matter),  $m_1$  ( $S_1$  field) and  $m_h$  (Higgs), plus the mixing angle  $\theta$  between  $S_1$  and  $h$ . We will fix the Higgs mass to  $m_h = 138$ GeV and give, in this section, the mixing angle  $\theta$  the two values  $10^\circ$  (small) and  $40^\circ$  (larger). The two last parameters we choose are the two physical mutual coupling constants  $\lambda_0^{(4)}$  (dark matter – Higgs) and  $\eta_{01}^{(4)}$  (dark matter –  $S_1$  particle), see (2.12).

In the framework of the thermal dynamics of the Universe within the standard cosmological model [32], the WIMP relic density is related to its annihilation rate by the familiar relations:

$$\begin{aligned} \Omega_D \bar{h}^2 &\simeq \frac{1.07 \times 10^9 x_f}{\sqrt{g_*} m_{\text{Pl}} \langle v_{12} \sigma_{\text{ann}} \rangle \text{GeV}}; \\ x_f &\simeq \ln \frac{0.038 m_{\text{Pl}} m_0 \langle v_{12} \sigma_{\text{ann}} \rangle}{\sqrt{g_*} x_f}. \end{aligned} \quad (3.1)$$

The notation is as follows: the quantity  $\bar{h}$  is the Hubble constant in units of  $100\text{km}/(\text{s} \times \text{Mpc})$ ,  $m_{\text{Pl}} = 1.22 \times 10^{19}\text{GeV}$  the Planck mass,  $m_0$  the dark matter mass,  $x_f = m_0/T_f$  the ratio of the dark matter mass to the freeze-out temperature  $T_f$  and  $g_*$  the number of relativistic degrees of freedom with a mass less than  $T_f$ . The quantity  $\langle v_{12} \sigma_{\text{ann}} \rangle$  is the thermally averaged annihilation cross-section of a pair of two dark matter particles multiplied by their relative speed in the center-of-mass reference frame. Solving (3.1) with the current value for the dark matter relic density  $\Omega_D \bar{h}^2 = 0.105 \pm 0.008$  [33] gives:

$$\langle v_{12} \sigma_{\text{ann}} \rangle \simeq (1.9 \pm 0.2) \times 10^{-9} \text{GeV}^{-2}, \quad (3.2)$$

for a range of dark matter masses between roughly 10GeV to 100GeV and  $x_f$  between 19.2 and 21.6, with about 0.4 thickness [34].

The value in (3.2) for the dark matter annihilation cross-section translates into a relation between the parameters of a given theory entering the calculated expression of  $\langle v_{12} \sigma_{\text{ann}} \rangle$ , hence imposing a constraint on these parameters which will limit the intervals of possible dark matter masses. This constraint can be exploited to examine aspects of the theory like perturbativity. For example, in our model, we can obtain via (3.2) the

mutual coupling constant  $\eta_{01}^{(4)}$  for given values of  $\lambda_0^{(4)}$ , study its behavior as a function of  $m_0$  and tell which dark-matter mass regions are consistent with perturbativity. Note that once the two mutual coupling constants  $\lambda_0^{(4)}$  and  $\eta_{01}^{(4)}$  are perturbative, all the other physical coupling constants will be. In the study of this section, we choose the values  $\lambda_0^{(4)} = 0.01$  (very weak), 0.2 (weak) and 1 (large). We also let the two masses  $m_0$  and  $m_1$  stretch from 0.1GeV to 120GeV, occasionally  $m_0$  to 200GeV. Finally, note that we do not incorporate the uncertainty in (3.2) when imposing the relic-density constraint, something that is sufficient in view of the descriptive nature of this work.

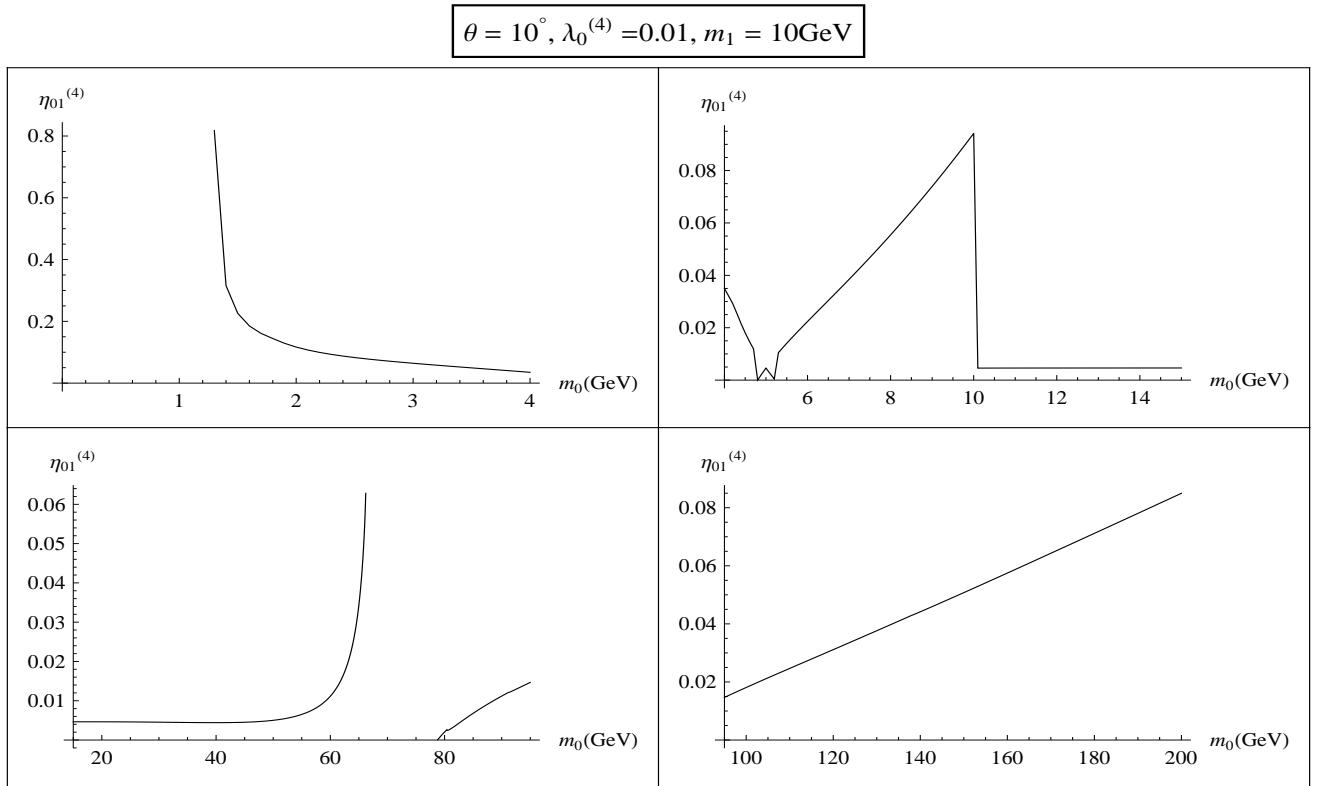


FIG. 1:  $\eta_{01}^{(4)}$  vs  $m_0$  for small  $m_1$ , small mixing and very small WIMP-Higgs coupling.

The dark matter annihilation cross sections (times the relative speed) through all possible channels are given in the appendix. The quantity  $\langle v_{12}\sigma_{\text{ann}} \rangle$  is the sum of all these contributions. Imposing  $\langle v_{12}\sigma_{\text{ann}} \rangle = 1.9 \times 10^{-9} \text{GeV}^{-2}$  dictates the behavior of  $\eta_{01}^{(4)}$ , which is displayed as a function of the dark matter mass  $m_0$ . Of course, as the parameters are numerous, the behavior is bound to be rich and diverse. We cannot describe every bit of it. Also, one has to note from the outset that for a given set of values for the parameters, the solution to the relic-density constraint is not unique: besides positive real solutions (when they exist), we may find negative real or even complex solutions. It is beyond the scope of the present work to investigate the nature and behavior of all the solutions. We

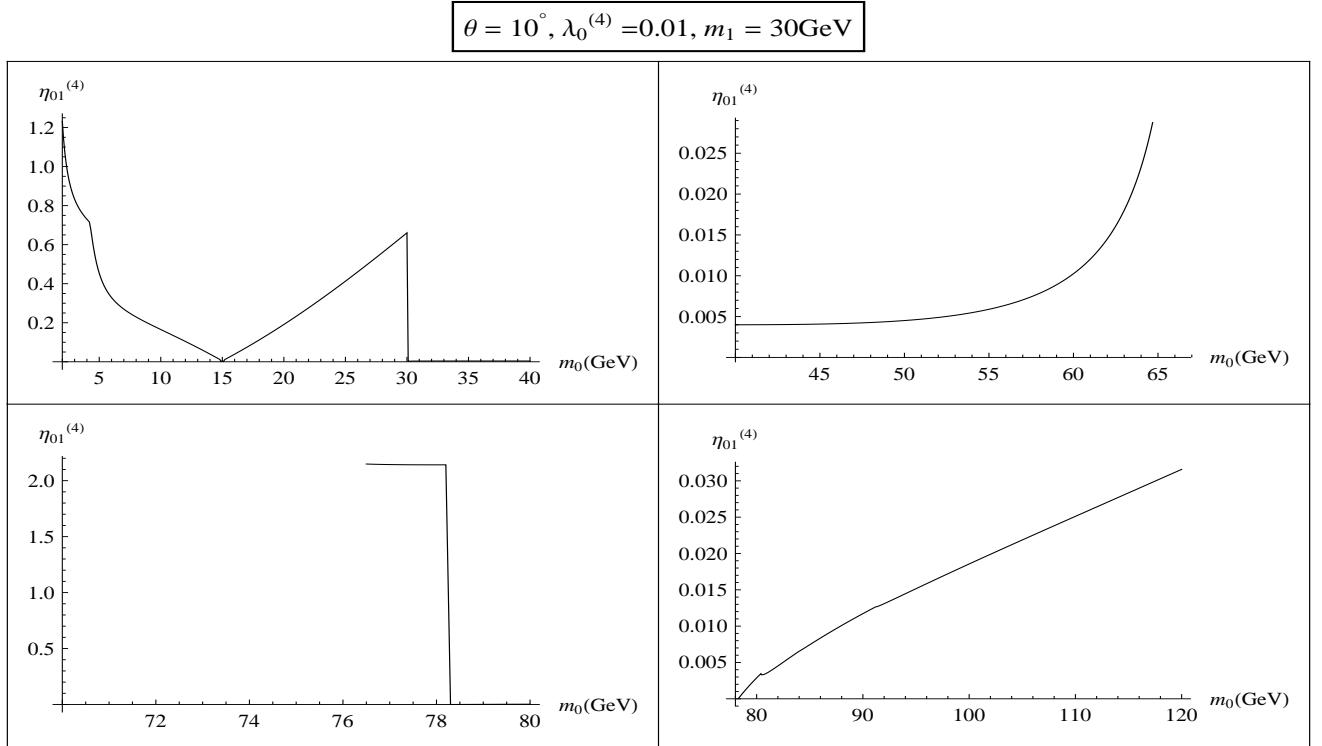


FIG. 2:  $\eta_{01}^{(4)}$  vs  $m_0$  for moderate  $m_1$ , small mixing and very small WIMP-Higgs coupling.

are only interested in finding the smallest positive real solution  $\eta_{01}^{(4)}$  when it exists, looking at its behavior and finding out when it is small enough to be perturbative.

### A. Small mixing angle and very weak dark matter – Higgs coupling

Let us describe briefly and only partly how the mutual  $S_0 - S_1$  coupling constant  $\eta_{01}^{(4)}$  behaves as a function of the  $S_0$  mass  $m_0$ . We start by a small mixing angle, say  $\theta = 10^\circ$ , and a very weak mutual  $S_0 -$  Higgs coupling constant, say  $\lambda_0^{(4)} = 0.01$ . Let us also fix the  $S_1$  mass first at the small value  $m_1 = 10\text{GeV}$ . The corresponding behavior of  $\eta_{01}^{(4)}$  versus  $m_0$  is shown in Fig. 1.

The range of  $m_0$  shown is from 0.1GeV to 200GeV, cut in four intervals to allow for ‘local’ features to be displayed\*. We see that the relic-density constraint on  $S_0$  annihilation has no positive real solution for  $m_0 \lesssim 1.3\text{GeV}$ , and so, with these very small masses,  $S_0$  cannot be a dark matter candidate. In other words, for  $m_1 = 10\text{GeV}$ , the particle  $S_0$  cannot annihilate into the lightest fermions only; inclusion of the  $c$ -quark is necessary. Note that right about  $m_0 \simeq 1.3\text{GeV}$ , the  $c$  threshold, the mutual coupling constant  $\eta_{01}^{(4)}$

\* A logplot in this descriptive study is not advisable.

$$\theta = 10^\circ, \lambda_0^{(4)} = 0.01, m_0 = 0.2 \text{ GeV}$$

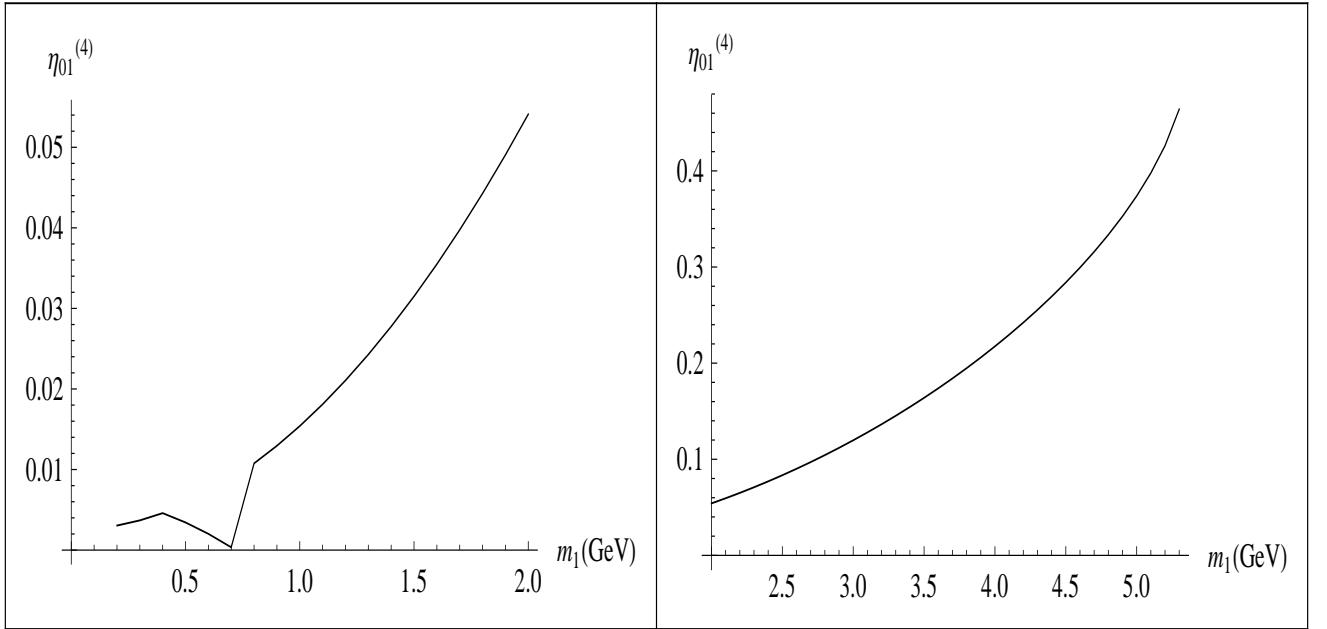


FIG. 3:  $\eta_{01}^{(4)}$  vs  $m_1$  for very light  $S_1$ , small mixing and very small WIMP-Higgs coupling.

starts at about 0.8, a value, while perturbative, that is roughly eighty-two-fold larger than the mutual  $S_0$  – Higgs coupling constant  $\lambda_0^{(4)}$ . Then  $\eta_{01}^{(4)}$  decreases, steeply first, more slowly as we cross the  $\tau$  mass towards the  $b$  mass. Just before  $m_1/2$ , the coupling  $\eta_{01}^{(4)}$  hops onto another solution branch that is just emerging from negative territory, gets back to the first one at precisely  $m_1/2$  as this latter carries now smaller values, and then jumps up again onto the second branch as the first crosses the  $m_0$ -axis down. It goes up this branch with a moderate slope until  $m_0$  becomes equal to  $m_1$ , a value at which the  $S_1$  annihilation channel opens. Right beyond  $m_1$ , there is a sudden fall to a value  $\eta_{01}^{(4)} \simeq 0.0046$  that is about half the value of  $\lambda_0^{(4)}$ , and  $\eta_{01}^{(4)}$  stays flat till  $m_0 \simeq 45 \text{ GeV}$  where it starts increasing, sharply after 60 GeV. In the mass interval  $m_0 \simeq 66 \text{ GeV} – 79 \text{ GeV}$ , there is a desert with no positive real solutions to the relic-density constraint, hence no viable dark matter candidate. Beyond  $m_0 \simeq 79 \text{ GeV}$ , the mutual coupling constant  $\eta_{01}^{(4)}$  keeps increasing monotonously, with a small notch at the  $W$  mass and a less noticeable one at the  $Z$  mass. Note that for this value of  $m_1$  (10 GeV), all values reached by  $\eta_{01}^{(4)}$  in the mass range considered, however large or small with respect to  $\lambda_0^{(4)}$ , are perturbatively acceptable.

Increasing  $m_1$  to moderate values does not change the above qualitative features. As an illustration, Fig. 2 shows the behavior of  $\eta_{01}^{(4)}$  as a function of  $m_0$  for  $m_1 = 30 \text{ GeV}$ , keeping the mixing angle  $\theta = 10^\circ$ , still small, and the mutual  $S_0$  – Higgs coupling constant

$\theta = 10^\circ, \lambda_0^{(4)} = 0.01, m_0 = 1.4\text{GeV(L)} \text{ and } 1.5\text{GeV(R)}$

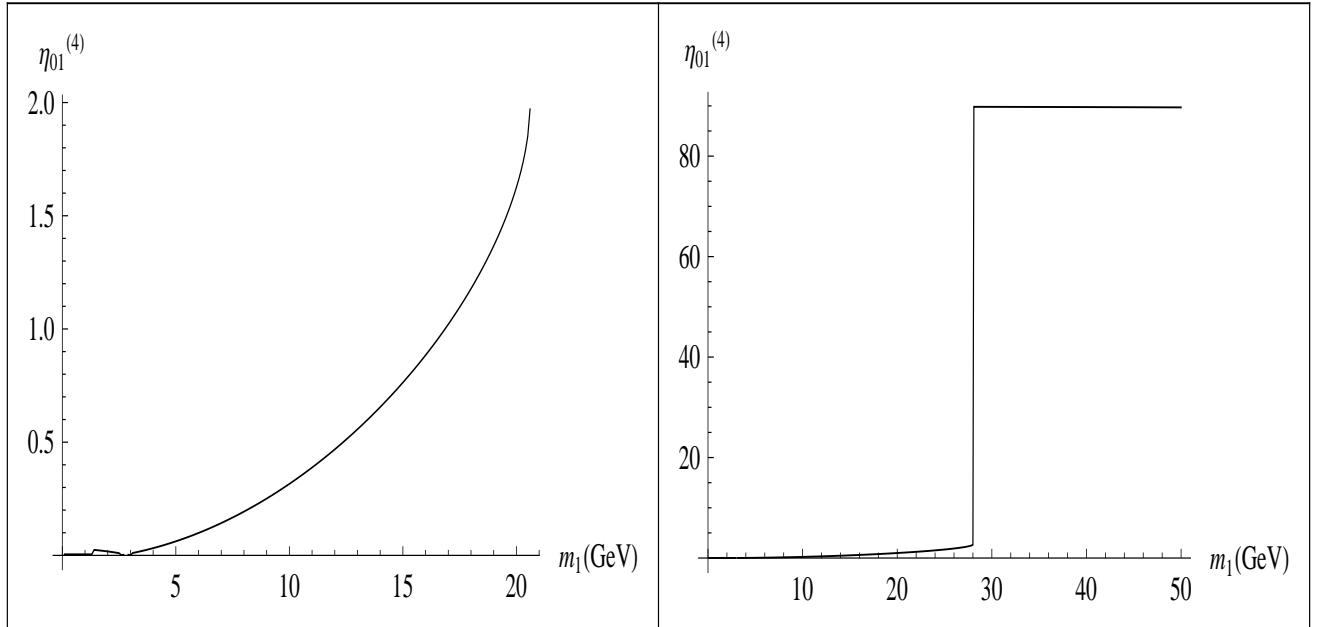


FIG. 4:  $\eta_{01}^{(4)}$  versus  $m_1$  for  $m_0$  above  $\tau$  threshold.

$\lambda_0^{(4)} = 0.01$ , still very weak. The first thing to note is that not all values of  $\eta_{01}^{(4)}$  are perturbative. Indeed,  $\eta_{01}^{(4)}$  does not start until  $m_0 \simeq 1.5\text{GeV}$ , but with the very large value<sup>†</sup> 89.8. It decreases very sharply right after, to 2.04 at about 1.6GeV. It continues to decrease with a pronounced change in the slope at the  $b$  threshold. Effects at the masses  $m_1/2$  and  $m_1$  similar to those of figure 1 do occur here too. There is a desert that lies in this case in the mass interval 66.5GeV – 76.5GeV. At the upper bound, the coupling  $\eta_{01}^{(4)}$  takes the value 2.15 and decreases very slowly till  $m_0 \simeq 78.2\text{GeV}$ . Right after this mass, it plunges down to catch up with a solution branch that is just emerging from negative values. This solution branch increases steadily with two small notches at the  $W$  and  $Z$  masses. A similar global behavior occurs at other moderate  $m_1$  masses, with varying local features.

Because of the very-small- $m_0$  deserts described and visible on Fig. 1, one may ask whether the model ever allows for very light dark matter. To look into this, we fix  $m_0$  at small values and let  $m_1$  vary. Take first  $m_0 = 0.2\text{GeV}$  and see Fig. 3. The allowed  $S_0$  annihilation channels are the very light fermions  $e, u, d, \mu$  and  $s$ , plus  $S_1$  when  $m_1 < m_0$ .

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<sup>†</sup> This feature is not displayed in figure 2 to avoid masking the other much smaller values taken by the mutual coupling.

$$\theta = 10^\circ, \lambda_0^{(4)} = 0.2, m_1 = 20 \text{ GeV}$$

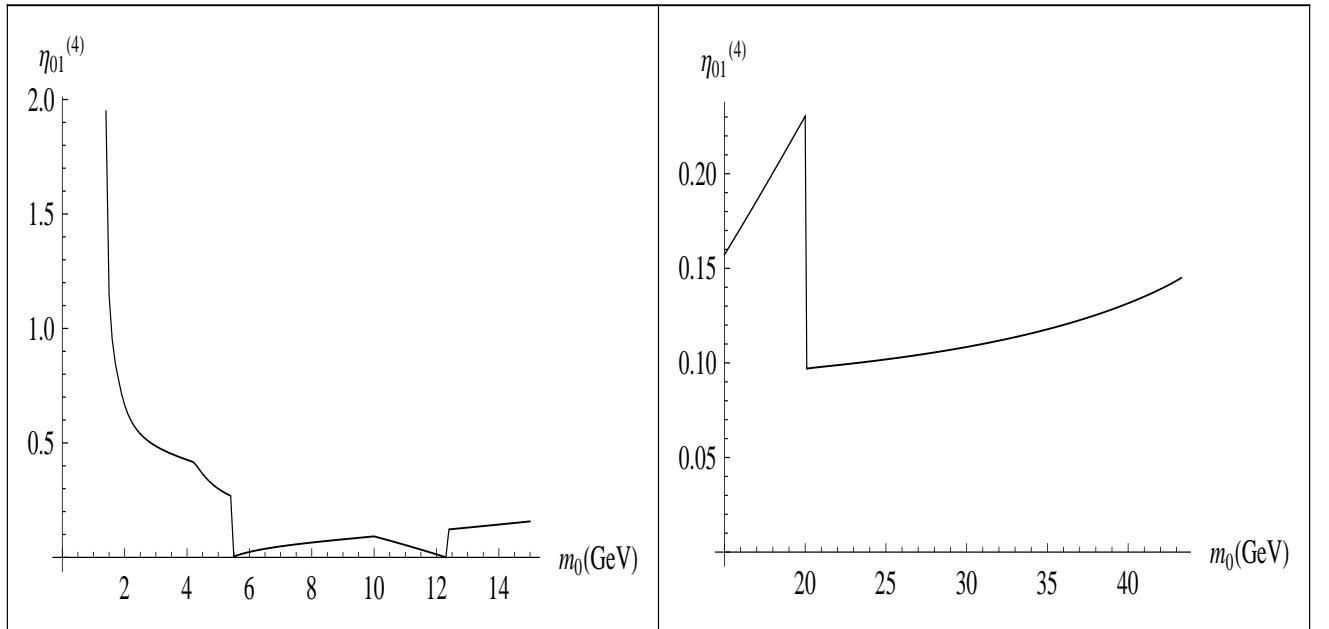


FIG. 5:  $\eta_{01}^{(4)}$  vs  $m_0$  for small mixing, moderate  $m_1$  and WIMP-Higgs coupling.

Note that we still have  $\theta = 10^\circ$  and  $\lambda_0^{(4)} = 0.01$ . Qualitatively, we notice that in fact, there are no solutions for  $m_1 < 0.2 \text{ GeV}$  ( $= m_0$  here), a mass at which  $\eta_{01}^{(4)}$  takes the very small value  $\simeq 0.003$ . It goes up a solution branch and leaves it at  $m_1 \simeq 0.4 \text{ GeV}$  to descend on a second branch that enters negative territory at  $m_1 \simeq 0.7 \text{ GeV}$ , forcing  $\eta_{01}^{(4)}$  to return onto the first branch. There is an accelerated increase till  $m_1 \simeq 5 \text{ GeV}$ , a value at which  $\eta_{01}^{(4)} \simeq 0.5$ . And then a desert, no positive real solutions, no viable dark matter.

Increasing  $m_0$  until about  $1.3 \text{ GeV}$  does not change these overall features: some ‘movement’ for very small values of  $m_1$  and then an accelerated increase till reaching a desert with a lower bound that changes with  $m_0$ . For example, the desert starts at  $m_1 \simeq 6.8 \text{ GeV}$  for  $m_0 = 0.6 \text{ GeV}$  and  $m_1 \simeq 7.3 \text{ GeV}$  for  $m_0 = 1.2 \text{ GeV}$ . Note that in all these cases where  $m_0 \lesssim 1.3 \text{ GeV}$ , all values of  $\eta_{01}^{(4)}$  are perturbative. Therefore, the model can very well accommodate very light dark matter with a restricted range of  $S_1$  masses.

However, the situation changes after the inclusion of the  $\tau$  annihilation channel. Indeed, as Fig. 4 shows, for  $m_0 = 1.4 \text{ GeV}$ , though the overall shape of the behavior of  $\eta_{01}^{(4)}$  as a function of  $m_1$  is qualitatively the same, the desert threshold is pushed significantly higher, to  $m_1 \simeq 20 \text{ GeV}$ . But more significantly,  $\eta_{01}^{(4)}$  starts to be larger than one already at  $m_1 \simeq 17 \text{ GeV}$ , therefore loosing perturbativity. For  $m_0 = 1.5 \text{ GeV}$ , the desert is effectively erased as we have a sudden jump to highly non-perturbative values of  $\eta_{01}^{(4)}$  right after  $m_1 \simeq 28 \text{ GeV}$ . Such a behavior stays with larger values of  $m_0$ . But for  $m_1 \lesssim 20 \text{ GeV}$  (case

$m_0 = 1.5\text{GeV}$ ), the values of  $\eta_{01}^{(4)}$  are smaller than one and physical use of the model is possible if needed.

$$\theta = 10^\circ, \lambda_0^{(4)} = 0.2, m_1 = 60\text{GeV}$$

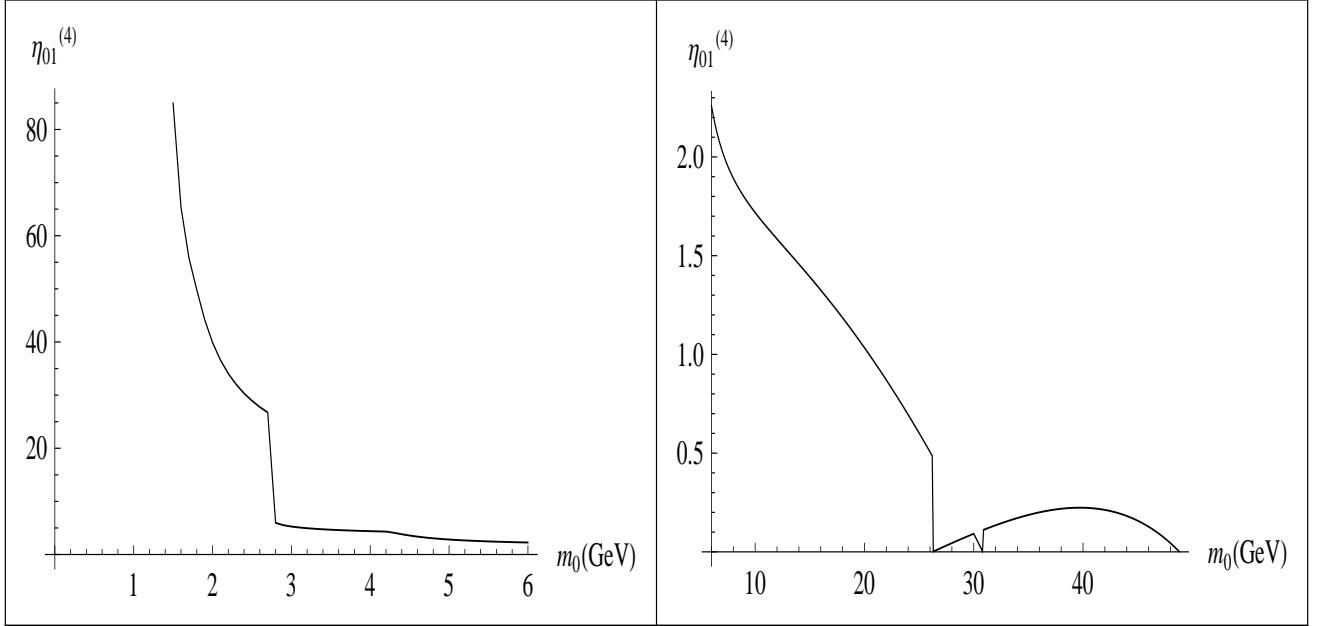


FIG. 6:  $\eta_{01}^{(4)}$  versus  $m_0$  for heavy  $S_1$ , small mixing and small WIMP-Higgs coupling.

### B. Small mixing angle and larger dark matter – Higgs couplings

What are the effects of the relic-density constraint when we vary the parameter  $\lambda_0^{(4)}$ ? Let us keep the Higgs –  $S_1$  mixing angle small ( $\theta = 10^\circ$ ) and increase  $\lambda_0^{(4)}$ , first to 0.2 and later to 1. For  $\lambda_0^{(4)} = 0.2$ , Figure 5 shows the behavior of  $\eta_{01}^{(4)}$  as a function of the dark matter mass  $m_0$  when  $m_1 = 20\text{GeV}$ . We see that  $\eta_{01}^{(4)}$  starts at  $m_0 \simeq 1.4\text{GeV}$  with a value of about 1.95. It decreases with a sharp change of slope at the  $b$  threshold, then makes a sudden dive at about 5 GeV, a change of branch at  $m_1/2$  down till about 12GeV where it jumps up back onto the previous branch just before going to cross into negative territory. It drops sharply at  $m_0 = m_1$  and then increases slowly until  $m_0 \simeq 43.3\text{GeV}$ . Beyond, there is nothing, a desert.

This is of course different from the situation of very small  $\lambda_0^{(4)}$  like in Fig. 1 and Fig. 2 above: here we see some kind of natural dark-matter mass ‘confinement’ to small-moderate

viable<sup>‡</sup> values.

Still for  $\lambda_0^{(4)} = 0.2$  with  $m_1 = 60\text{GeV}$  this time, Fig. 6 shows  $\eta_{01}^{(4)}$  starting very high ( $\simeq 85\text{GeV}$ ) at  $m_0 \simeq 1.5\text{GeV}$ , decreasing quickly with a first sudden drop at  $2.7\text{ GeV}$  and a second one to zero at  $26.2\text{ GeV}$ . A solution branch is then picked up – left briefly at  $m_1/2$  – until  $49\text{ GeV}$  and then nothing. What is peculiar here is that, in contrast with previous situations, the desert starts at a mass  $m_0 < m_1$ , i.e., before the opening of the  $S_1$  annihilation channel. In other words, the dark matter is annihilating into the light fermions only and the model is perturbatively viable in the range  $20\text{GeV} – 49\text{GeV}$ .

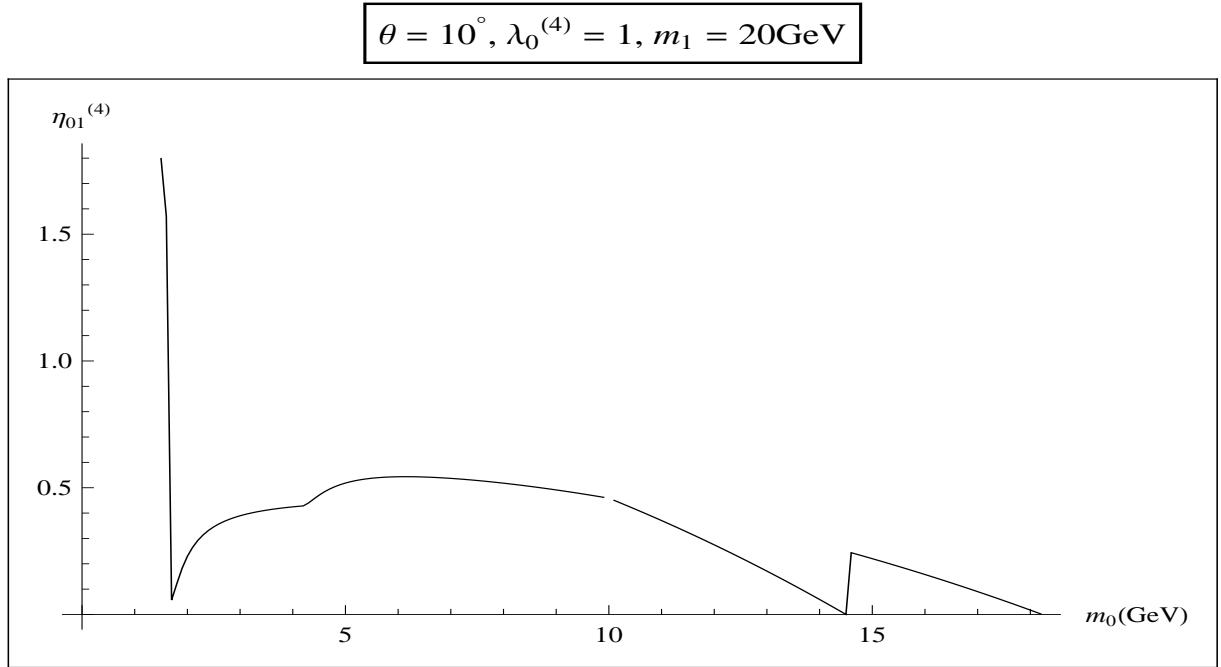


FIG. 7:  $\eta_{01}^{(4)}$  versus  $m_0$  for medium  $m_1$ , small mixing and large WIMP-higgs coupling.

The case  $\lambda_0^{(4)} = 1$  with  $m_1 = 20\text{GeV}$  is displayed in Fig. 7. There are no solutions below  $m_0 \simeq 1.5\text{GeV}$  at which  $\eta_{01}^{(4)} \simeq 1.80$ . From this value,  $\eta_{01}^{(4)}$  slips down very quickly to pick up less abruptly when crossing the  $\tau$  threshold. There is a significant change in the slope at the crossing of the  $b$  mass. Note the absence of a solution at  $m_1/2$ , which is a new feature, present for other values of  $m_1$  not displayed here. Beyond  $m_1/2$ , there is a slight change in the downward slope, a change of solution branch, and that goes until  $14.5\text{GeV}$  where  $\eta_{01}^{(4)}$  jumps to catch up with the previous branch. It goes down this branch until about  $18\text{GeV}$  where the desert starts.

We have studied the behavior of  $\eta_{01}^{(4)}$  as a function of  $m_0$  for other values of  $m_1$  between  $20\text{GeV}$  and  $100\text{GeV}$  while keeping  $\theta = 10^\circ$  and  $\lambda_0^{(4)} = 1$ . For  $m_1 \lesssim 79.2\text{GeV}$ , the behavior

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<sup>‡</sup> Note that the values of  $\eta_{01}^{(4)}$  for  $1.6\text{GeV} \lesssim m_0 \lesssim 43.3\text{GeV}$  are all perturbative.

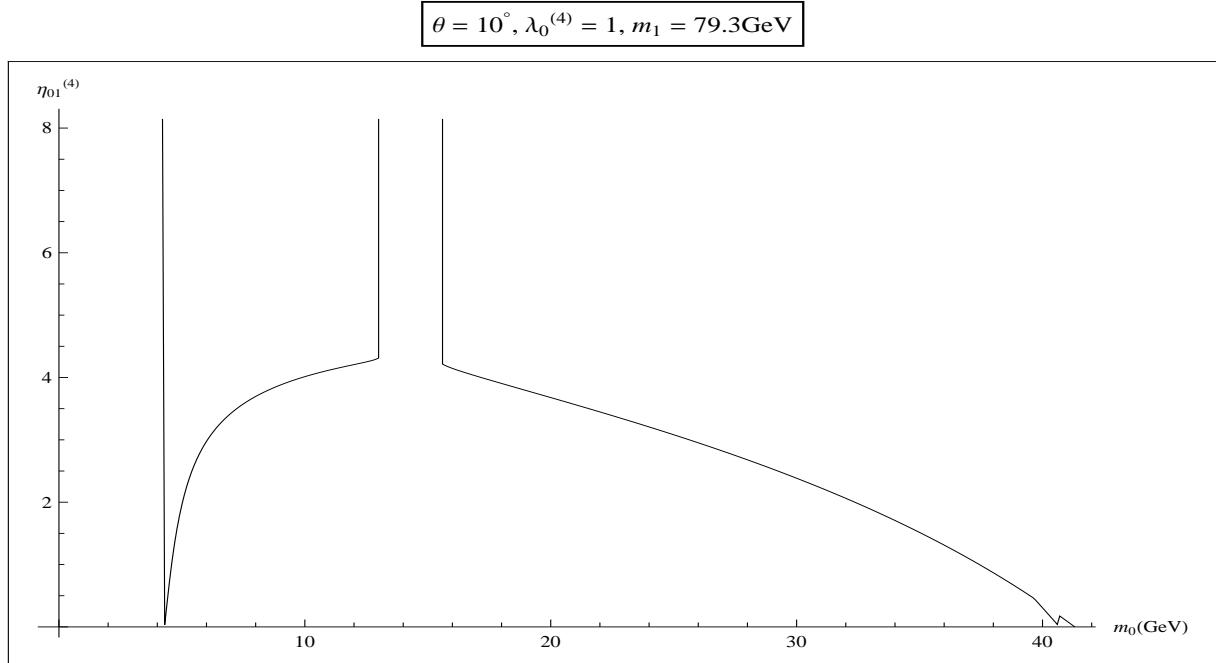


FIG. 8:  $\eta_{01}^{(4)}$  versus  $m_0$  for heavy  $S_1$ , small mixing and large WIMP-Higgs coupling.

is qualitatively quite similar to that shown in Fig. 7, but beyond this mass, there is a highly non-perturbative branch  $\eta_{01}^{(4)}$  jumps onto at small and moderate values of  $m_0$ . This highly non-perturbative region stretches in size as  $m_1$  increases. Fig. 8 displays this new feature. Note that on this figure, not all of the range of  $\eta_{01}^{(4)}$  is shown in order to allow the small-coupling regions to be displayed; the high values of  $\eta_{01}^{(4)}$  are in the two thousands. Note also that it is the same highly non-perturbative solution branch  $\eta_{01}^{(4)}$  jumps onto for other large values of  $m_1$ .

### C. Larger mixing angles

Last in this descriptive study is to see the effects of larger values of the  $S_1$  – Higgs mixing angle  $\theta$ . We give it here the value  $\theta = 40^\circ$  and tune back the mutual  $S_0$  – Higgs coupling constant  $\lambda_0^{(4)}$  to the very small value 0.01. Figure 9 shows the behavior of  $\eta_{01}^{(4)}$  as a function of  $m_0$  for  $m_1 = 20\text{GeV}$ . One recognizes features similar to those of the case  $\theta = 10^\circ$ , though coming in different relative sizes. The very-small- $m_0$  desert ends at about 0.3GeV. There are by-now familiar features at the  $c$  and  $b$  masses,  $m_1/2$  and  $m_1$ . Two relatively small forbidden intervals (deserts) appear for relatively large values of the dark matter mass: 67.3GeV – 70.9GeV and 79.4GeV – 90.8GeV. The  $W$  mass region is forbidden but there is action as we cross the  $Z$  mass.

Other values of  $m_1$ , not displayed because of space, behave similarly with an additional

$$\theta = 40^\circ, \lambda_0^{(4)} = 0.01, m_1 = 20\text{GeV}$$

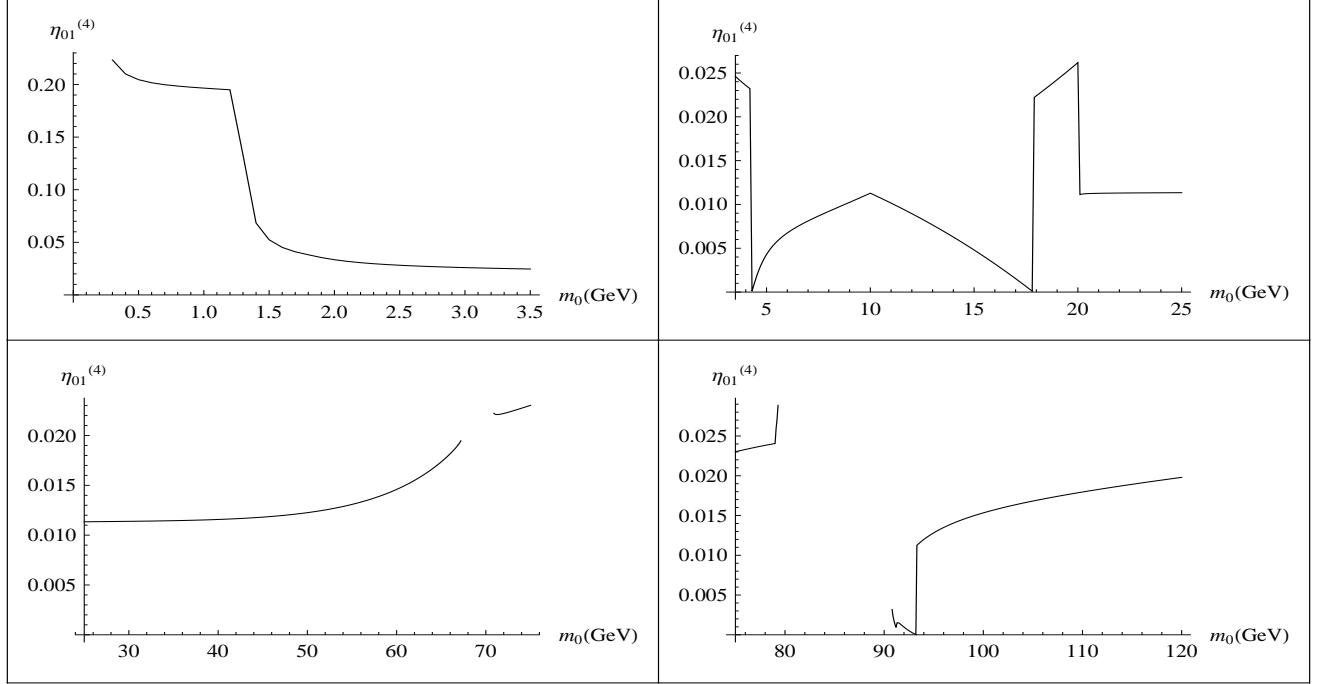


FIG. 9:  $\eta_{01}^{(4)}$  versus  $m_0$  for moderate  $m_1$ , moderate mixing and small WIMP-Higgs coupling.

effect, namely, a sudden drop in slope at  $m_0 = (m_h + m_1)/2$  coming from the ignition of  $S_0$  annihilation into  $S_1$  and Higgs. We have also worked out the cases  $\lambda_0^{(4)} = 0.2$  and 1 for  $\theta = 40^\circ$ . The case  $\lambda_0^{(4)} = 0.2$  is displayed in Fig. 10 and presents differences with the corresponding small-mixing situation  $\theta = 10^\circ$ . Indeed, for  $m_1 = 20\text{GeV}$ , the first feature we notice is a smoother behavior; compare with Fig. 5. Here,  $\eta_{01}^{(4)}$  starts at  $m_0 \simeq 0.3\text{GeV}$  with the small value  $\simeq 0.016$  and goes up, faster at the  $c$  mass and with a small effect at the  $b$  mass. It increases very slowly until  $m_1/2$  and decreases very slowly until  $m_0 = m_1$ , and then there is a sudden change of branch followed immediately by a desert<sup>§</sup>. So here too the model naturally confines the mass of a viable dark matter to small-moderate values, a dark matter particle annihilating into light fermions only. What is also noticeable is that there is stability of  $\eta_{01}^{(4)}$  around the value of  $\lambda_0^{(4)}$  in the interval  $1.5\text{GeV} - 20\text{GeV}$  ( $= m_1$  here).

The case  $m_1 = 60\text{GeV}$  presents also overall similarities as well as noticeable differences with the corresponding case  $\theta = 10^\circ$ , see Fig. 6. The first difference is that all values of  $\eta_{01}^{(4)}$  are perturbative. This latter starts at  $m_0 \simeq 1.4\text{GeV}$  with the value  $\sim 0.75$ , goes down and jumps to catch up with another solution branch emerging from negative territory when

<sup>§</sup> Except for the very tiny interval  $78.5\text{GeV} - 79.0\text{GeV}$  not displayed on Fig. 10.

$$\theta = 40^\circ, \lambda_0^{(4)} = 0.2, m_1 = 20\text{GeV(L)} \text{ and } 60\text{GeV(R)}$$

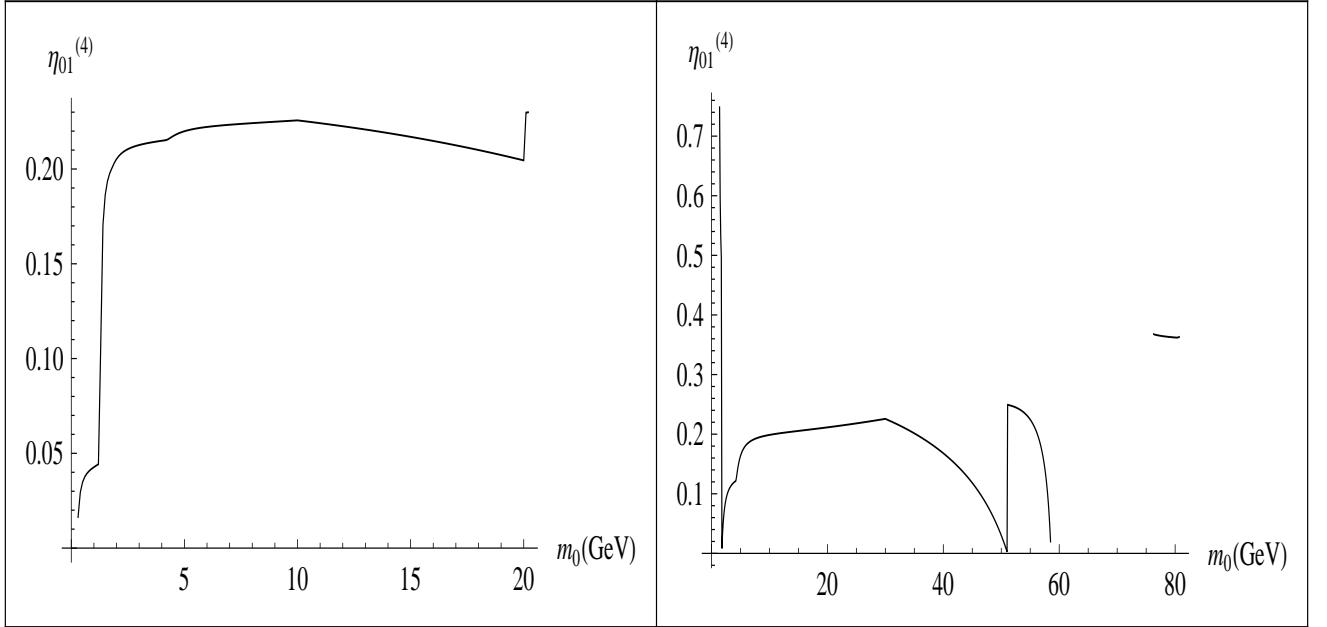


FIG. 10:  $\eta_{01}^{(4)}$  versus  $m_0$  for moderate (L) and large (R)  $m_1$ , large mixing and moderate WIMP-Higgs coupling.

crossing the  $\tau$  mass. It increases, kicking up when crossing the  $b$ -quark mass. It changes slope down at  $m_1/2$  and goes to zero at about 51GeV. It jumps up onto another branch that goes down to zero also at about 58.6GeV, just below  $m_1$ , and then there is a desert, except for the small interval 76.3GeV – 80.5GeV.

The case  $\lambda_0^{(4)} = 1$  is shown in Fig. 11. Global similarities with the previous case are apparent. All values of  $\eta_{01}^{(4)}$  are perturbative and the mass range is naturally confined to the interval 0.2GeV – 20GeV for  $m_1 = 20\text{GeV}$ , and 1.4GeV – 52.3GeV for  $m_1 = 60\text{GeV}$ . We note action at the usual masses and, in particular, we see there are no solutions at  $m_0 = m_1/2$  like in the case  $\theta = 10^\circ$ . We note here too the quasi-constancy of  $\eta_{01}^{(4)}$  for most of the available range.

Finally, we note that we have worked out larger mixing angles, notably  $\theta = 75^\circ$ . In general, these cases do not display any new features worth discussing: the overall behavior is mostly similar to what we have seen, with expected relative variations in size.

$$\theta = 40^\circ, \lambda_0^{(4)} = 1, m_1 = 20\text{GeV(L)} \text{ and } 60\text{GeV(R)}$$

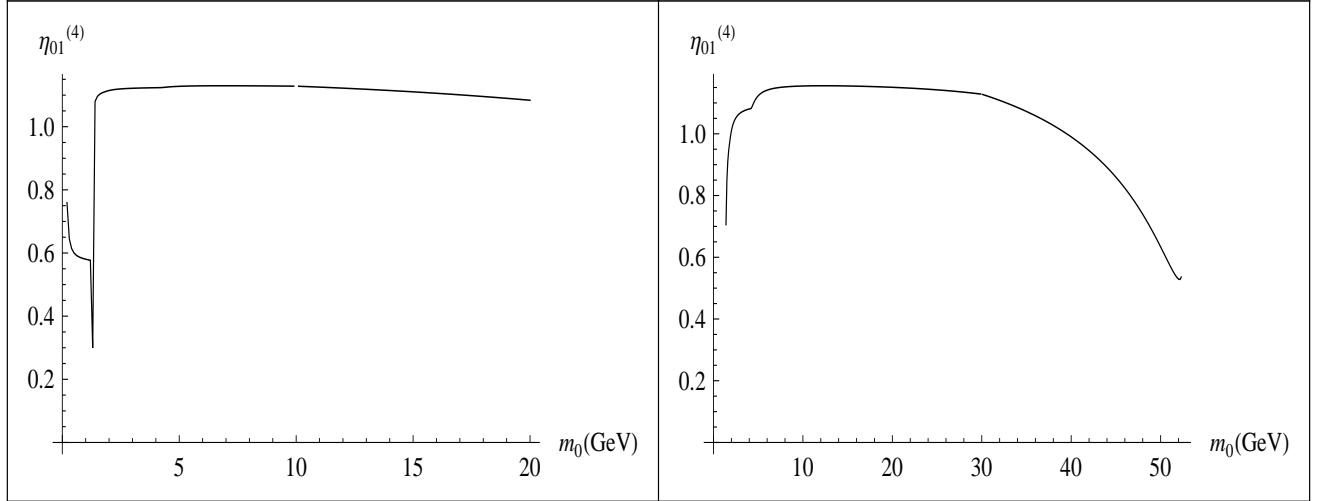


FIG. 11:  $\eta_{01}^{(4)}$  versus  $m_0$  for moderate (L) and large (R)  $m_1$ , large mixing and large WIMP-Higgs coupling.

#### IV. DARK-MATTER DIRECT DETECTION

Experiments like CDMS II [27? ], XENON 10/100 [8, 26], DAMA/LIBRA [2] and CoGeNT [3] search directly for a dark matter signal. Such a signal would typically come from the elastic scattering of a dark matter WIMP off a non-relativistic nucleon target. However, throughout the years, such experiments have not yet detected an unambiguous signal, but rather yielded increasingly stringent exclusion bounds on the dark matter – nucleon elastic-scattering total cross section  $\sigma_{\text{det}}$  in terms of the dark matter mass  $m_0$ .

In order for a theoretical dark-matter model to be viable, it has to satisfy these bounds. It is therefore natural to inquire whether the model we present in this work has any capacity of describing dark matter. Hence, we have to calculate  $\sigma_{\text{det}}$  as a function of  $m_0$  for different values of the parameters  $(\theta, \lambda_0^{(4)}, m_1)$  and project its behavior against the experimental bounds. We will limit ourselves to the region  $0.1\text{GeV} - 100\text{GeV}$  as we are interested in light dark matter. As experimental bounds, we will use the results from CDMSII and XENON100, as well as the future projections of SuperCDMS [28] and XENON1T [29]. The results of CoGeNT, DAMA/LIBRA and CRESST will be discussed elsewhere [35]. As the figures below show [36], in the region of our interest, XENON100 is only slightly tighter than CDMSII, SuperCDMS significantly lower and XENON1T the most stringent by far. But it is important to note that all these results loose reasonable predictability in the very light sector, say below 5GeV.

The scattering of  $S_0$  off a SM fermion  $f$  occurs via the t-channel exchange of the SM higgs and  $S_1$ . In the non-relativistic limit, the effective Lagrangian describing this

interaction reads

$$\mathcal{L}_{S_0-f}^{(eff)} = a_f \bar{f} f S_0^2, \quad (4.1)$$

where

$$a_f = -\frac{m_f}{2v} \left[ \frac{\lambda_0^{(3)} \cos \theta}{m_h^2} - \frac{\eta_{01}^{(3)} \sin \theta}{m_1^2} \right]. \quad (4.2)$$

In this case the total cross section for this process is given by :

$$\sigma_{S_0 f \rightarrow S_0 f} = \frac{m_f^4}{4\pi (m_f + m_0)^2 v^2} \left[ \frac{\lambda_0^{(3)} \cos \theta}{m_h^2} - \frac{\eta_{01}^{(3)} \sin \theta}{m_1^2} \right]^2. \quad (4.3)$$

At the nucleon level, the effective interaction between a nucleon  $N = p$  or  $n$  and  $S_0$  has the form

$$\mathcal{L}_{S_0-N}^{(eff)} = a_N \bar{N} N S_0^2, \quad (4.4)$$

where the effective nucleon- $S_0$  coupling constants is given by:

$$a_N = \frac{(m_N - \frac{7}{9}m_B)}{v} \left[ \frac{\lambda_0^{(3)} \cos \theta}{m_h^2} - \frac{\eta_{01}^{(3)} \sin \theta}{m_1^2} \right]. \quad (4.5)$$

In this relation,  $m_N$  is the nucleon mass and  $m_B$  the baryon mass in the chiral limit [25]. The total cross section for non-relativistic  $S_0 - N$  elastic scattering is therefore:

$$\sigma_{\text{det}} \equiv \sigma_{S_0 N \rightarrow S_0 N} = \frac{m_N^2 (m_N - \frac{7}{9}m_B)^2}{4\pi (m_N + m_0)^2 v^2} \left[ \frac{\lambda_0^{(3)} \cos \theta}{m_h^2} - \frac{\eta_{01}^{(3)} \sin \theta}{m_1^2} \right]^2. \quad (4.6)$$

The rest of this section is devoted to a brief discussion of the behavior of  $\sigma_{\text{det}}$  as a function of  $m_0$ . We will of course impose the relic-density constraint on the dark matter annihilation cross section (3.2). But in addition, we will require that the coupling constants are perturbative, and so impose the additional requirement  $0 \leq \eta_{01}^{(4)} \leq 1$ . Also, here too, the choices of the sets of values of the parameters  $(\theta, \lambda_0^{(4)}, m_1)$  can by no means be exhaustive but only indicative. Furthermore, though a detailed description of the behavior of  $\sigma_{\text{det}}$  could be interesting in its own right, we will refrain from doing so in this work as there is no need for it, and content ourselves with mentioning overall features and trends. Generally, as  $m_0$  increases, the detection cross section  $\sigma_{\text{det}}$  starts from high values, slopes down to minima that depend on the parameters and then picks up moderately. There are features and action at the usual mass thresholds, with varying sizes and shapes. Excluded regions are there, those coming from the relic-density constraint and new ones originating from the additional perturbativity requirement. Close to the upper boundary of the mass interval considered in this study, there is no universal behavior to mention as in some cases

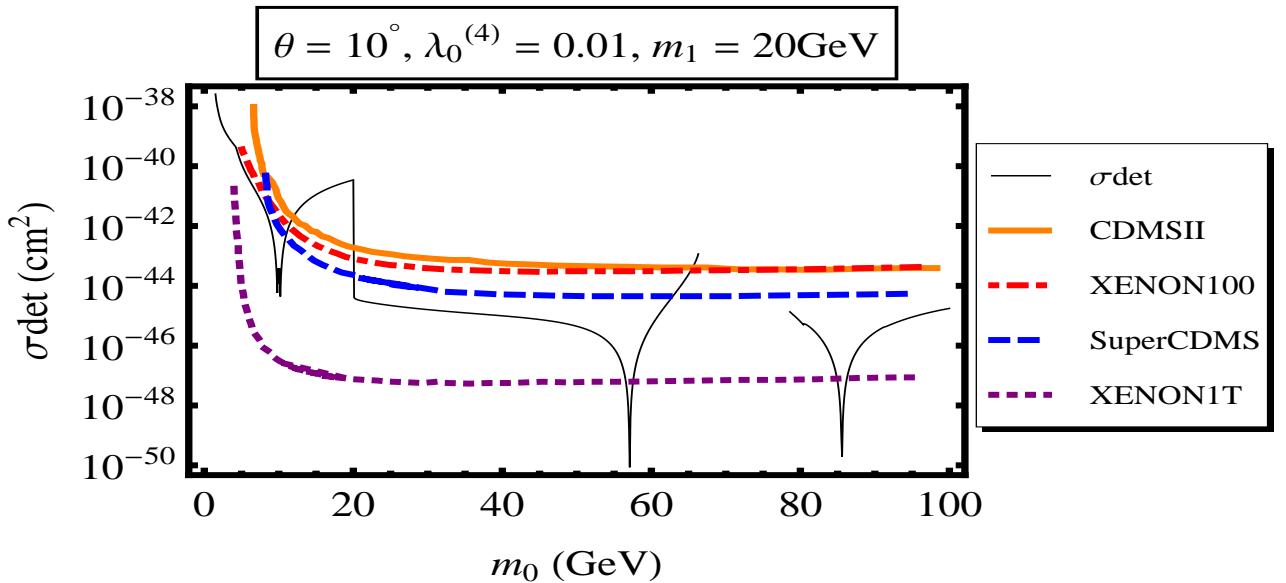


FIG. 12: *Elastic  $N - S_0$  scattering cross-section as a function of  $m_0$  for moderate  $m_1$ , small mixing and small WIMP-Higgs coupling.*

$\sigma_{\text{det}}$  will increase monotonously and, in some others, it will decrease or ‘not be there’ at all. Let us finally remark that the logplots below may not show these general features clearly as these latter are generally distorted.

Let us start with the small Higgs –  $S_1$  mixing angle  $\theta = 10^\circ$  and the very weak mutual  $S_0$  – Higgs coupling  $\lambda_0^{(4)} = 0.01$ . Fig. 12 shows the behavior of  $\sigma_{\text{det}}$  versus  $m_0$  in the case  $m_1 = 20\text{GeV}$ . We see that for the two mass intervals  $20\text{GeV} - 65\text{GeV}$  and  $75\text{GeV} - 100\text{GeV}$ , plus an almost singled-out peaks at  $m_0 = m_1/2$ , the elastic scattering cross section is below the projected sensitivity of SuperCDMS. However, XENON1T will probe all the these masses , except  $m_0 \simeq 58$  GeV and 85 GeV.

Also, as we see in Fig. 12, most of the mass range for very light dark matter is excluded for these values of the parameters. Is this systematic? In general, smaller values of  $m_1$  drive the predictability ranges to the lighter sector of the dark matter masses. Figure 13 illustrates this pattern. We have taken  $m_1 = 5\text{GeV}$ , just above the lighter quarks threshold. In the small-mass region, we see that SuperCDMS is passed in the range  $5\text{GeV} - 30\text{GeV}$ . Again , all this mass ranges will be probed by XENON1T experiment, except a sharp peak at  $m_0 = m_1/2 = 2.5\text{GeV}$ , but for such a very light mass, the experimental results are not without ambiguity.

Reversely, increasing  $m_1$  shuts down possibilities for very light dark matter and thins the intervals as it drives the predicted masses to larger values. For instance, in figure 14 where  $m_1 = 40\text{GeV}$ , in addition to the peak at  $m_1/2$  that crosses SuperCDMS but not XENON1T, we see acceptable masses in the ranges  $40\text{GeV} - 65\text{GeV}$  and  $78\text{GeV} - \text{up}$ .

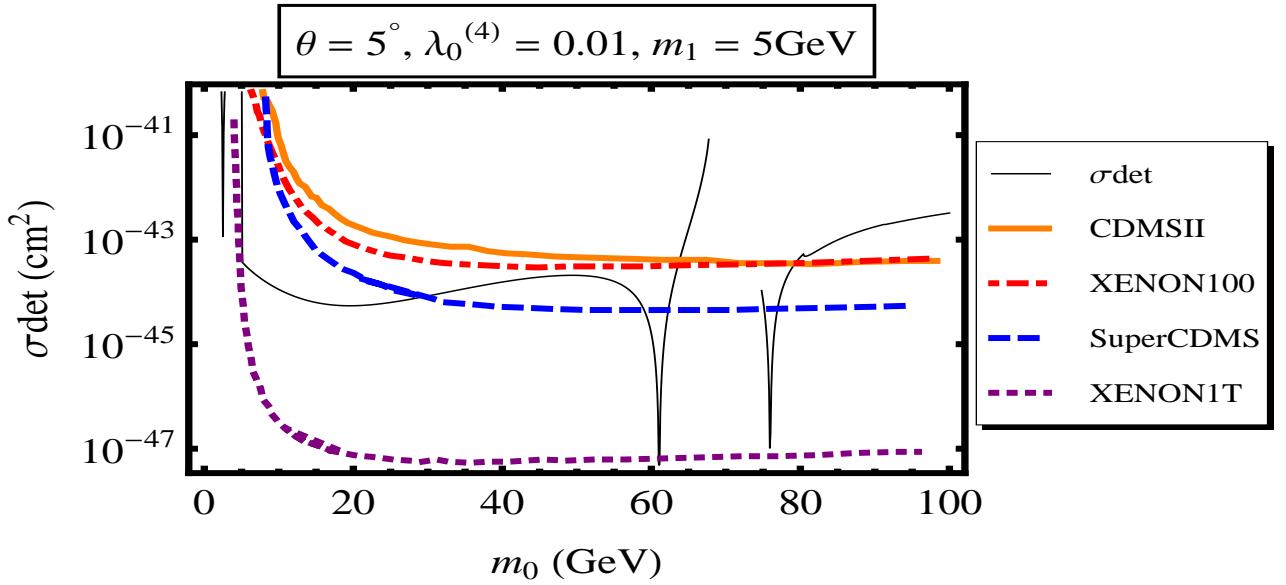


FIG. 13: *Elastic  $N - S_0$  scattering cross-section as a function of  $m_0$  for light  $S_1$ , small mixing and small WIMP-Higgs coupling.*

Here too the intervals narrow as we descend, surviving XENON1T as spiked peaks at 62 GeV and around 95 GeV.

A larger mutual coupling constant  $\lambda_0^{(4)}$  has the general effect of squeezing the acceptable intervals of  $m_0$  by pushing the values of  $\sigma_{\text{det}}$  up. As an illustration, see figure 15 where we have taken  $\lambda_0^{(4)} = 0.2$  and a larger value of  $m_1 = 60\text{GeV}$ . In this example, already XENON100 excludes all the masses below 100 GeV except a relatively narrow peak at  $m_1/2$ .

Increasing the mixing angle  $\theta$  has also the general effect of increasing the value of  $\sigma_{\text{det}}$ . Figure 16 shows this trend for  $\theta = 40^\circ$ ; compare with Fig. 12. The only allowed masses by the current bounds of CDMSII and XENON100 are between 20 GeV and 50 GeV, the narrow interval around  $m_1/2$ , and another very sharp one, at about 94GeV. The projected sensitivity of XENON1T will probe all mass ranges except those at  $m_0 \simeq 30\text{GeV}$  and 94GeV.

Finally, it happens that there are regions of the parameters for which the model has no predictability. See figure 17 for illustration. We have combined the effects of increasing the values of the two parameters  $\lambda_0^{(4)}$  and  $m_1$ . As we see, we barely get something at  $m_1/2$  that cannot even cross XENON100 down to SuperCDMS.

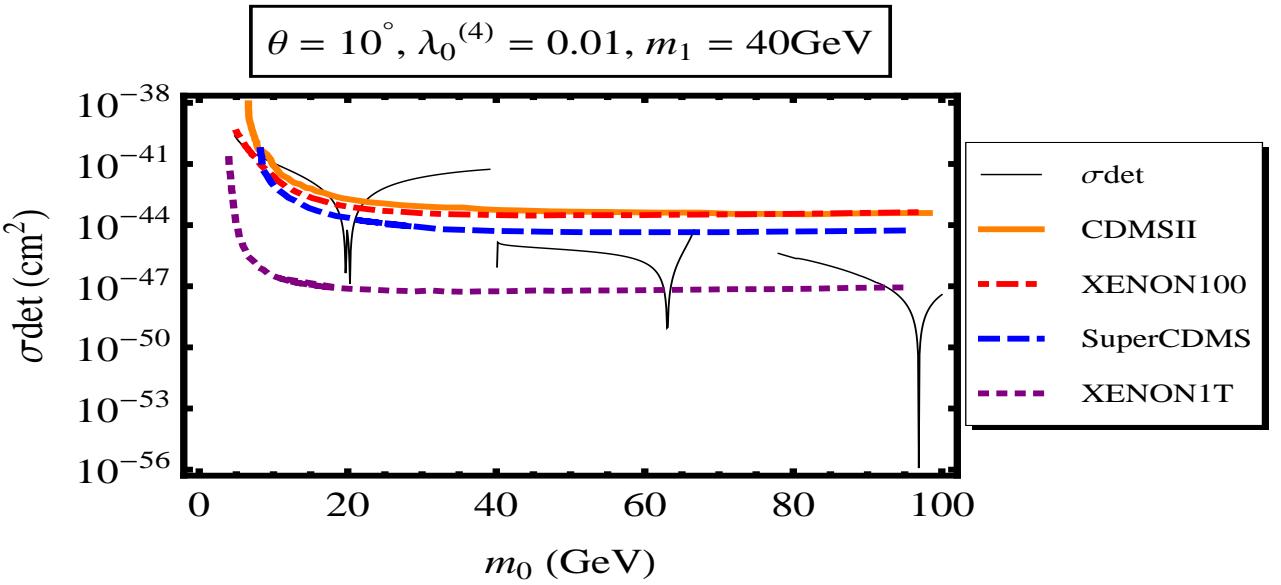


FIG. 14: *Elastic  $N - S_0$  scattering cross-section as a function of  $m_0$  for medium  $m_1$ , small mixing and small WIMP-Higgs coupling.*

## V. CONCLUDING REMARKS

In this work, we presented a plausible scenario for light cold dark matter, (for masses lighter than 100 GeV). This latter consists in enlarging the Standard Model with two gauge-singlet  $\mathbb{Z}_2$ -symmetric scalar fields. One is the dark matter field  $S_0$ , stable, while the other undergoes spontaneous symmetry breaking, resulting in the physical field  $S_1$ . This opens additional channels through which  $S_0$  can annihilate, hence a reducing its number density. The model is parametrized by three quantities: the physical mutual coupling constant  $\lambda_0^{(4)}$  between  $S_0$  and the Higgs, the mixing angle  $\theta$  between  $S_1$  and the Higgs and the mass  $m_1$  of the particle  $S_1$ . We first imposed on  $S_0$  annihilation cross section the constraint from the observed dark-matter relic density and studied its effects through the behavior of the physical mutual coupling constant  $\eta_{01}^{(4)}$  between  $S_0$  and  $S_1$  as a function of the dark matter mass  $m_0$ . Apart from forbidden regions (deserts) or others where perturbativity is lost, we find that for most values of the three parameters, there are viable solutions in the small-moderate masses of the dark matter sector. Deserts are found for most of the ranges of the parameters whereas perturbativity is lost mainly for larger values of  $m_1$ . Through the behavior of  $\eta_{01}^{(4)}$ , we could see the mass thresholds which mostly affect the annihilation of dark matter, and these are at the  $c$ ,  $\tau$  and  $b$  masses, as well as  $m_1/2$  and  $m_1$ .

The current experimental bounds from CDMSII and XENON100 put a strong constraint on the  $S_0$  masses in the range between 10 to 20 GeV. For small values of  $m_1$ , very light dark matter is viable, with a mass as small as one GeV. This is of course useful for

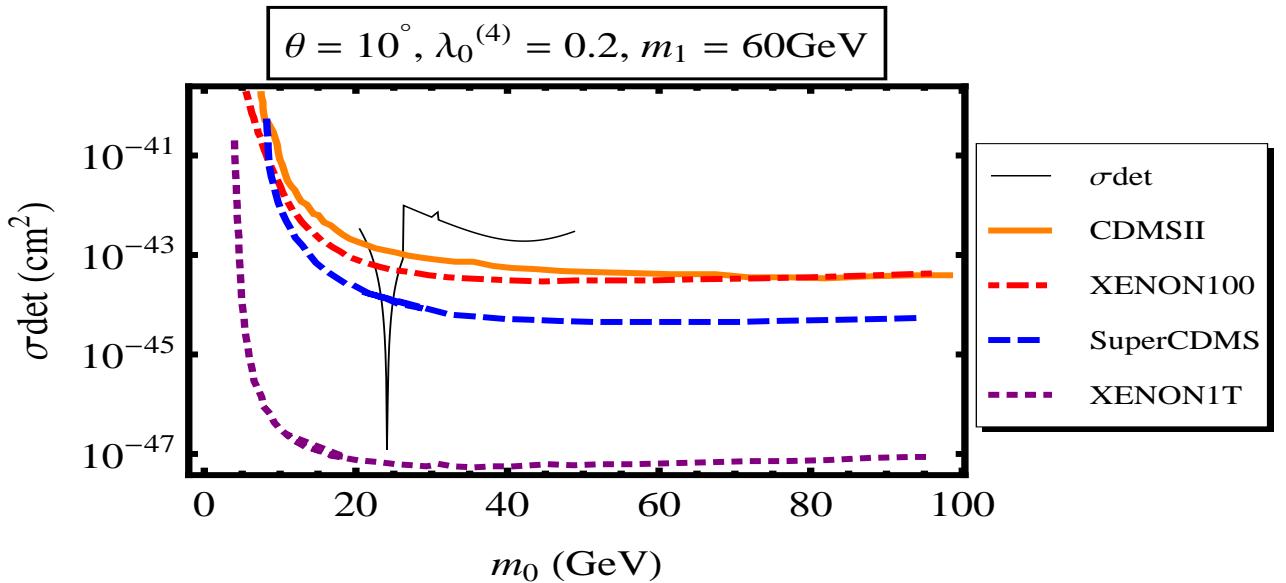


FIG. 15: *Elastic  $N - S_0$  scattering cross-section as a function of  $m_0$  for heavy  $S_1$ , small mixing and moderate WIMP-Higgs coupling.*

understanding the results of the experiments DAMA/LIBRA, CoGeNT , CRESST [14] as well as the recent data of the Fermi Gamma Ray Space Telescope [4] . The projected sensitivity of future WIMP direct searches such as XENON1T will probe all the  $S_0$  masses between 5 GeV and 100 GeV.

The next step to take is to test the model against the phenomenological constraints. Indeed, one important feature of the model is that it mixes the  $S_1$  field with the Higgs. This must have implications on the Higgs detection through the measurable channels. Current experimental bounds from LEP II data can be used to constrain our mixing angle  $\theta$ , and possibly other parameters. In addition, a very light  $S_0$  and/or  $S_1$  will contribute to the invisible decay of  $J/\psi$  and  $\Upsilon$  mesons and can lead to a significant branching fraction. These constraints can be injected back into the model and restrain further its domain of validity. These issues are under current investigation [35].

Also, in this work, the  $S_1$  vacuum expectation value  $v_1$  was taken equal to 100GeV, but a priori, nothing prevents us from considering other scales. However, taking  $v_1$  much larger than the electro-weak scale requires  $\eta_{01}^{(4)}$  to be very tiny , which will result in the suppression of the crucial annihilation channel  $S_0S_0 \rightarrow S_1S_1$  . Also, we have fixed the Higgs mass to  $m_h = 138\text{GeV}$ , which is consistent with the current acceptable experimental bounds [30]. Yet, it can be useful to ask here too what the effect of changing this mass would be.

Finally, in this study, besides the dark matter field  $S_0$ , only one extra field has been considered. Naturally, one can generalize the investigation to include  $N$  such fields and

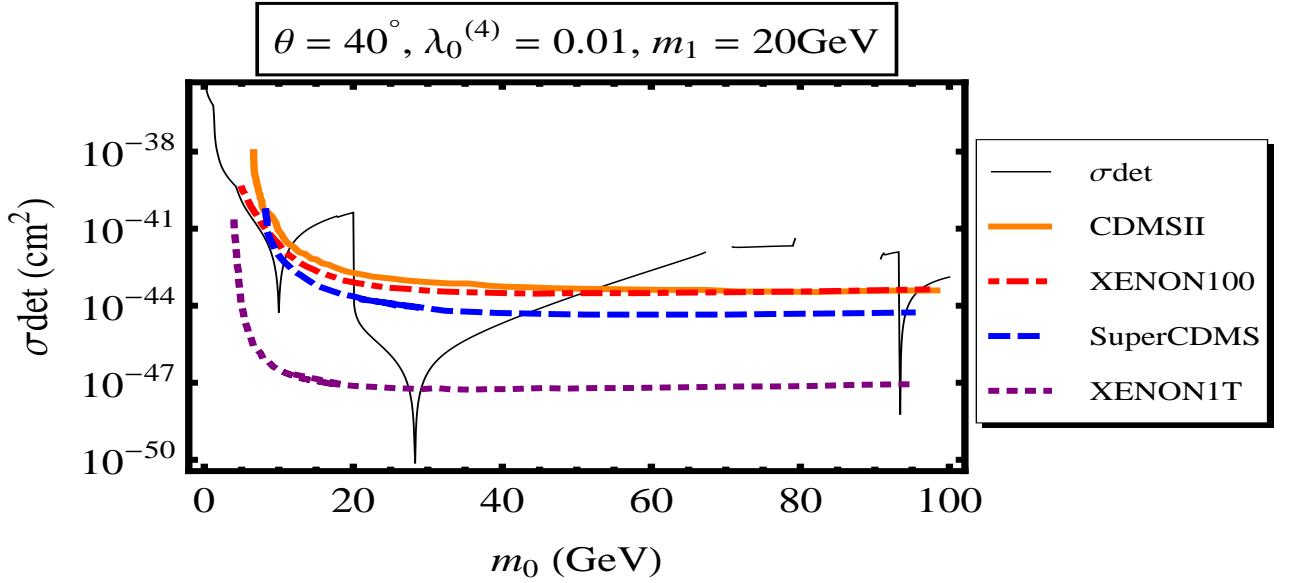


FIG. 16: *Elastic  $N - S_0$  scattering cross-section as a function of  $m_0$  for moderate  $m_1$ , large mixing and small WIMP-Higgs coupling.*

discuss the cosmology and particle phenomenology in terms of  $N$ . It just happens that the model is rich enough to open new possibilities in the quest of dark matter worth pursuing.

### Appendix A: Dark matter annihilation cross-sections

The cross-sections related to the annihilation  $S_0$  into the scalar particles are as follows. For the  $hh$  channel, we have:

$$v_{12}\sigma_{S_0S_0 \rightarrow hh} = \frac{\sqrt{m_0^2 - m_h^2}}{64\pi m_0^3} \Theta(m_0 - m_h) \left[ \left( \lambda_0^{(4)} \right)^2 + \frac{4\lambda_0^{(4)} \left( \lambda_0^{(3)} \right)^2}{m_h^2 - 2m_0^2} + \frac{2\lambda_0^{(4)} \lambda_0^{(3)} \lambda^{(3)}}{4m_0^2 - m_h^2} \right. \\ + \frac{2\lambda_0^{(4)} \lambda_1^{(3)} \eta_{01}^{(3)} (4m_0^2 - m_1^2)}{(4m_0^2 - m_1^2)^2 + \epsilon_1^2} + \frac{4 \left( \lambda_0^{(3)} \right)^4}{(m_h^2 - 2m_0^2)^2} + \frac{4\lambda^{(3)} \left( \lambda_0^{(3)} \right)^3}{(4m_0^2 - m_h^2)(m_h^2 - 2m_0^2)} \\ + \frac{4 \left( \lambda_0^{(3)} \right)^2 \lambda_1^{(3)} \eta_{01}^{(3)} (4m_0^2 - m_1^2)}{[(4m_0^2 - m_1^2)^2 + \epsilon_1^2] (m_h^2 - 2m_0^2)} + \frac{(\lambda^{(3)})^2 \left( \lambda_0^{(3)} \right)^2}{(4m_0^2 - m_h^2)^2} \\ \left. + \frac{\left( \lambda_1^{(3)} \right)^2 \left( \eta_{01}^{(3)} \right)^2}{(4m_0^2 - m_1^2)^2 + \epsilon_1^2} + \frac{2\lambda_0^{(3)} \lambda_1^{(3)} \lambda^{(3)} \eta_{01}^{(3)} (4m_0^2 - m_1^2)}{[(4m_0^2 - m_1^2)^2 + \epsilon_1^2] (4m_0^2 - m_h^2)} \right]. \quad (A1)$$

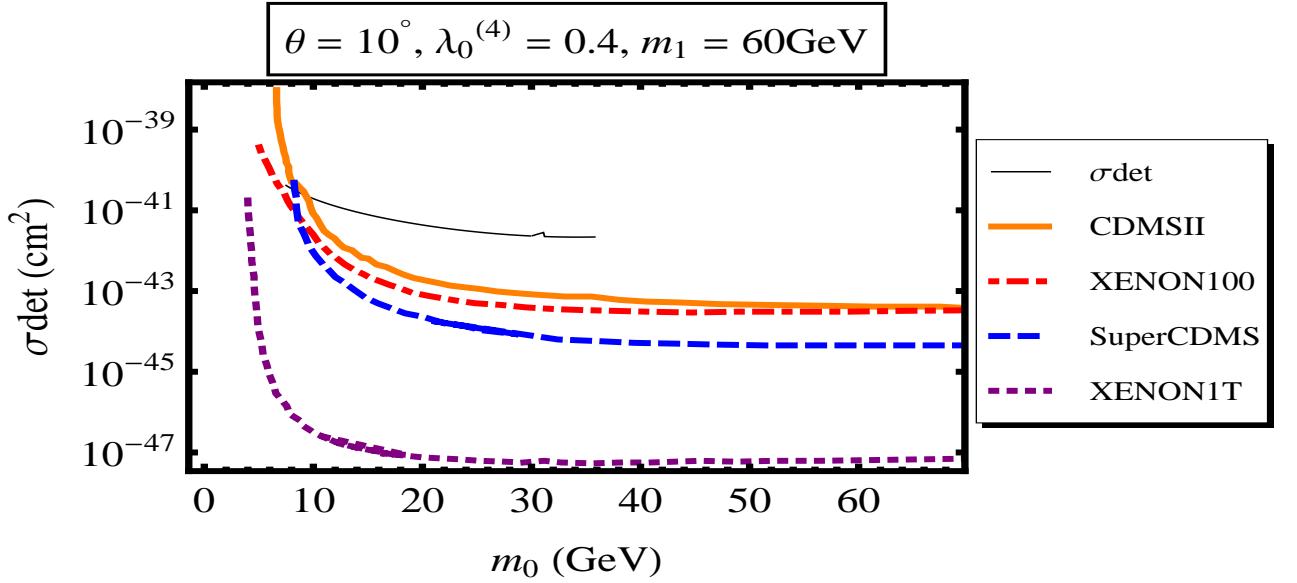


FIG. 17: elastic cross-section  $\sigma_{el}$  as a function of  $S_0$  mass for heavy  $S_1$ , small mixing and relatively large WIMP-Higgs coupling.

The  $\Theta$  function is the step function. For the  $S_1 S_1$  channel, we have the result:

$$\begin{aligned}
 v_{12} \sigma_{S_0 S_0 \rightarrow S_1 S_1} = & \frac{\sqrt{m_0^2 - m_1^2}}{64\pi m_0^3} \Theta(m_0 - m_1) \left[ \left( \eta_{01}^{(4)} \right)^2 + \frac{4\eta_{01}^{(4)} \left( \eta_{01}^{(3)} \right)^2}{m_1^2 - 2m_0^2} + \frac{2\eta_{01}^{(4)} \eta_{01}^{(3)} \eta_1^{(3)}}{4m_0^2 - m_1^2} \right. \\
 & + \frac{2\eta_{01}^{(4)} \lambda_0^{(3)} \lambda_2^{(3)} (4m_0^2 - m_h^2)}{(4m_0^2 - m_h^2)^2 + \epsilon_h^2} + \frac{4 \left( \eta_{01}^{(3)} \right)^4}{(m_1^2 - 2m_0^2)^2} + \frac{4 \left( \eta_{01}^{(3)} \right)^3 \eta_1^{(3)}}{(4m_0^2 - m_1^2)(m_1^2 - 2m_0^2)} \\
 & + \frac{4 \left( \eta_{01}^{(3)} \right)^2 \lambda_0^{(3)} \lambda_2^{(3)} (4m_0^2 - m_h^2)}{[(4m_0^2 - m_h^2)^2 + \epsilon_h^2] (m_1^2 - 2m_0^2)} + \frac{\left( \eta_{01}^{(3)} \right)^2 \left( \eta_1^{(3)} \right)^2}{(4m_0^2 - m_1^2)^2} \\
 & \left. + \frac{\left( \lambda_0^{(3)} \right)^2 \left( \lambda_2^{(3)} \right)^2}{(4m_0^2 - m_h^2)^2 + \epsilon_h^2} + \frac{2\eta_{01}^{(3)} \eta_1^{(3)} \lambda_0^{(3)} \lambda_2^{(3)} (4m_0^2 - m_h^2)}{[(4m_0^2 - m_h^2)^2 + \epsilon_h^2] (4m_0^2 - m_1^2)} \right]. \quad (A2)
 \end{aligned}$$

For the  $hS_1$  channel, we have:

$$\begin{aligned}
v_{12}\sigma_{S_0S_0 \rightarrow S_1h} = & \frac{\sqrt{[4m_0^2 - (m_h - m_1)^2][4m_0^2 - (m_h + m_1)^2]}}{128\pi m_0^4} \Theta(2m_0 - m_h - m_1) \\
& \left[ \left( \lambda_{01}^{(4)} \right)^2 + \frac{8\lambda_{01}^{(4)}\eta_{01}^{(3)}\lambda_0^{(3)}}{m_h^2 + m_1^2 - 4m_0^2} + \frac{2\lambda_{01}^{(4)}\lambda_0^{(3)}\lambda_1^{(3)}}{4m_0^2 - m_h^2} + \frac{2\lambda_{01}^{(4)}\eta_{01}^{(3)}\lambda_2^{(3)}}{4m_0^2 - m_1^2} \right. \\
& + \frac{16\left(\eta_{01}^{(3)}\right)^2\left(\lambda_0^{(3)}\right)^2}{(m_h^2 + m_1^2 - 4m_0^2)^2} + \frac{8\left(\lambda_0^{(3)}\right)^2\eta_{01}^{(3)}\lambda_1^{(3)}}{(m_h^2 + m_1^2 - 4m_0^2)(4m_0^2 - m_h^2)} \\
& + \frac{8\left(\eta_{01}^{(3)}\right)^2\lambda_0^{(3)}\lambda_2^{(3)}}{(m_h^2 + m_1^2 - 4m_0^2)(4m_0^2 - m_1^2)} + \frac{\left(\lambda_0^{(3)}\right)^2\left(\lambda_1^{(3)}\right)^2}{(4m_0^2 - m_h^2)^2} \\
& \left. + \frac{2\eta_{01}^{(3)}\lambda_0^{(3)}\lambda_1^{(3)}\lambda_2^{(3)}}{(4m_0^2 - m_h^2)(4m_0^2 - m_1^2)} + \frac{\left(\lambda_2^{(3)}\right)^2\left(\eta_{01}^{(3)}\right)^2}{(4m_0^2 - m_1^2)^2} \right]. \tag{A3}
\end{aligned}$$

The annihilation cross-section into fermions is:

$$\begin{aligned}
v_{12}\sigma_{S_0S_0 \rightarrow f\bar{f}} = & \frac{\sqrt{\left(m_0^2 - m_f^2\right)^3}}{4\pi m_0^3} \Theta(m_0 - m_f) \left[ \frac{\left(\lambda_0^{(3)}\lambda_{hf}\right)^2}{\left(4m_0^2 - m_h^2\right)^2 + \epsilon_h^2} + \frac{\left(\eta_{01}^{(3)}\lambda_{1f}\right)^2}{\left(4m_0^2 - m_1^2\right)^2 + \epsilon_1^2} \right. \\
& \left. + \frac{2\lambda_0^{(3)}\eta_{01}^{(3)}\lambda_{hf}\lambda_{1f}\left(4m_0^2 - m_h^2\right)\left(4m_0^2 - m_1^2\right)}{\left[\left(4m_0^2 - m_h^2\right)^2 + \epsilon_h^2\right]\left[\left(4m_0^2 - m_1^2\right)^2 + \epsilon_1^2\right]} \right]. \tag{A4}
\end{aligned}$$

The annihilation cross-section into  $W$ 's is given by:

$$\begin{aligned}
v_{12}\sigma_{S_0S_0 \rightarrow WW} = & \frac{\sqrt{m_0^2 - m_w^2}}{16\pi m_0^3} \Theta(m_0 - m_w) \left[ 1 + \frac{\left(2m_0^2 - m_w^2\right)^2}{2m_w^4} \right] \\
& \times \left[ \frac{\left(\lambda_0^{(3)}\lambda_{hw}\right)^2}{\left(4m_0^2 - m_h^2\right)^2 + \epsilon_h^2} + \frac{\left(\eta_{01}^{(3)}\lambda_{1w}\right)^2}{\left(4m_0^2 - m_1^2\right)^2 + \epsilon_1^2} \right. \\
& \left. + \frac{2\lambda_0^{(3)}\eta_{01}^{(3)}\lambda_{hw}\lambda_{1w}\left(4m_0^2 - m_h^2\right)\left(4m_0^2 - m_1^2\right)}{\left[\left(4m_0^2 - m_h^2\right)^2 + \epsilon_h^2\right]\left[\left(4m_0^2 - m_1^2\right)^2 + \epsilon_1^2\right]} \right]. \tag{A5}
\end{aligned}$$

Last, the annihilation cross-section into  $Z$ 's is:

$$\begin{aligned}
v_{12}\sigma_{S_0S_0 \rightarrow ZZ} &= \frac{\sqrt{m_0^2 - m_z^2}}{8\pi m_0^3} \Theta(m_0 - m_z) \left[ 1 + \frac{(2m_0^2 - m_z^2)^2}{2m_z^4} \right] \\
&\times \left[ \frac{\left( \lambda_0^{(3)} \lambda_{hz}^{(3)} \right)^2}{(4m_0^2 - m_h^2)^2 + \epsilon_h^2} + \frac{\left( \eta_{01}^{(3)} \lambda_{1z}^{(3)} \right)^2}{(4m_0^2 - m_1^2)^2 + \epsilon_1^2} \right. \\
&\left. + \frac{2\lambda_0^{(3)} \eta_{01}^{(3)} \lambda_{hz}^{(3)} \lambda_{1z}^{(3)} (4m_0^2 - m_h^2) (4m_0^2 - m_1^2)}{\left[ (4m_0^2 - m_h^2)^2 + \epsilon_h^2 \right] \left[ (4m_0^2 - m_1^2)^2 + \epsilon_1^2 \right]} \right]. \tag{A6}
\end{aligned}$$

The quantities  $\epsilon_h = m_h \Gamma_h$  and  $\epsilon_1 = m_1 \Gamma_1$  are regulators at the respective resonances. The decay rates  $\Gamma_h$  and  $\Gamma_1$  are calculable in perturbation theory. We have for  $h$ :

$$\begin{aligned}
\epsilon_{h \rightarrow f\bar{f}} &= \frac{(\lambda_{hf})^2}{8\pi} m_h^2 N_c \left( 1 - \frac{4m_f^2}{m_h^2} \right)^{\frac{3}{2}} \Theta(m_h - 2m_f); \\
\epsilon_{h \rightarrow WW} &= \frac{(\lambda_{hw}^{(3)})^2}{8\pi} \left( 1 - \frac{4m_w^2}{m_h^2} \right)^{\frac{1}{2}} \left[ 1 + \frac{(m_h^2 - 2m_w^2)^2}{8m_w^4} \right] \Theta(m_h - 2m_w); \\
\epsilon_{h \rightarrow ZZ} &= \frac{(\lambda_{hz}^{(3)})^2}{4\pi} \left( 1 - \frac{4m_z^2}{m_h^2} \right)^{\frac{1}{2}} \left[ 1 + \frac{(m_h^2 - 2m_z^2)^2}{8m_z^4} \right] \Theta(m_h - 2m_z); \\
\epsilon_{h \rightarrow S_0S_0} &= \frac{(\lambda_0^{(3)})^2}{32\pi} \left( 1 - \frac{4m_0^2}{m_h^2} \right)^{\frac{1}{2}} \Theta(m_h - 2m_0); \\
\epsilon_{h \rightarrow S_1S_1} &= \frac{(\lambda_2^{(3)})^2}{32\pi} \left( 1 - \frac{4m_1^2}{m_h^2} \right)^{\frac{1}{2}} \Theta(m_h - 2m_1). \tag{A7}
\end{aligned}$$

For  $S_1$ , we have similar expressions:

$$\begin{aligned}
\epsilon_{S_1 \rightarrow f\bar{f}} &= \frac{(\lambda_{1f})^2}{8\pi} m_1^2 N_c \left(1 - \frac{4m_f^2}{m_1^2}\right)^{\frac{3}{2}} \Theta(m_1 - 2m_f); \\
\epsilon_{S_1 \rightarrow WW} &= \frac{(\lambda_{1w}^{(3)})^2}{8\pi} \left(1 - \frac{4m_w^2}{m_1^2}\right)^{\frac{1}{2}} \left[1 + \frac{(m_1^2 - 2m_w^2)^2}{8m_w^4}\right] \Theta(m_1 - 2m_w); \\
\epsilon_{S_1 \rightarrow ZZ} &= \frac{(\lambda_{1z}^{(3)})^2}{4\pi} \left(1 - \frac{4m_z^2}{m_1^2}\right)^{\frac{1}{2}} \left[1 + \frac{(m_1^2 - 2m_z^2)^2}{8m_z^4}\right] \Theta(m_1 - 2m_z); \\
\epsilon_{S_1 \rightarrow S_0 S_0} &= \frac{(\eta_{01}^{(3)})^2}{32\pi} \left(1 - \frac{4m_0^2}{m_1^2}\right)^{\frac{1}{2}} \Theta(m_1 - 2m_0); \\
\epsilon_{S_1 \rightarrow hh} &= \frac{(\lambda_1^{(3)})^2}{32\pi} \left(1 - \frac{4m_h^2}{m_1^2}\right)^{\frac{1}{2}} \Theta(m_1 - 2m_h). \tag{A8}
\end{aligned}$$

where  $N_c$  is equal to 1 for leptons and 3 for quarks.

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